

Probing baryogenesis with gravitational waves

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Affleck-Dine baryogenesis is compelling yet challenging to probe because of the high-energy physics involved. We demonstrate that this mechanism can be generically realized with low-energy new physics without supersymmetry while producing detectable gravitational waves (GWs) sourced by parametric resonance of a light scalar field. In viable models, the scalar has a mass of $\mathcal{O}(0.1\text{--}10)$ GeV, yielding GWs with peak frequencies of $\mathcal{O}(10\text{--}100)$ Hz. This study further reveals a new complementarity between upcoming LIGO-frequency GW detectors and laboratory searches across frontiers of particle physics.

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Introduction. The observed excess of matter over antimatter in the universe, in terms of the baryon-to-photon ratio $\eta = \frac{n_b - n_{\bar{b}}}{n_\gamma} \approx 6 \times 10^{-10}$ [1], remains a long-standing puzzle in particle physics and cosmology. Affleck-Dine (AD) baryogenesis [2] is a compelling solution to the puzzle, where the asymmetry emerges from the oscillation and decay of a baryon number-carrying complex scalar field. Originally formulated with supersymmetry (SUSY) [2–5], it involves flat direction moduli fields composed of superpartners of standard model (SM) particles. These light moduli fields can acquire large displacements during inflation, break C and CP symmetries, and subsequently roll down and oscillate to produce baryon asymmetry, while the inflaton still dominates the Universe [6] (alternatively, the AD field itself can be the inflaton [7,8]).

Despite its appeal, experimentally testing the AD mechanism is challenging because of the high-scale physics involved in the inflationary era. Fortunately, the stochastic gravitational wave background (SGWB) generated by the

oscillating AD field offers a promising new probe by utilizing the naturally high-energy environment in the early cosmos: oscillating scalar fields can trigger *parametric resonance*, i.e., rapid nonperturbative particle production, leading to large inhomogeneities and fragmentation that source GWs (e.g., [9–17]).

In conventional SUSY AD baryogenesis, parametric resonance does not generate detectable GWs [18–21]. Meanwhile, SUSY AD models can lead to Q-balls, which may result in early matter domination and induce GWs at second order by enhancing preexisting curvature perturbations [18,19]. Q-ball-induced GWs strongly depend on the speed of the transition from matter to radiation domination [22].

The purpose of this work is twofold: first, we present non-SUSY realizations of the AD mechanism that involve new physics below the electroweak scale and can operate in cosmic epochs long after inflation; second, we demonstrate that parametric resonance sourced by AD condensates in such models naturally produces SGWB within reach of upcoming GW experiments.

This study uncovers a generic class of well-motivated sources of SGWB and adds a timely complement to the rapidly growing field of probing particle physics with GWs, building on the ground-breaking discoveries of GWs by LIGO-Virgo-KAGRA collaboration [23–25] and the recent strong evidence for SGWBs revealed by pulsar timing arrays [26–33]. Based on simple benchmark models, we show that successful baryogenesis requires new physics in the AD sector with a characteristic scale $m_\Phi = \mathcal{O}(0.1\text{--}10)$ GeV, and the resulting GWs have a peak frequency in the range of

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$\mathcal{O}(1\text{--}100)$ Hz, within reach of upcoming detectors, including the Cosmic Explorer (CE) [34] and Einstein Telescope (ET) [35], as well as DECIGO [36] and Big Bang Observatory (BBO) [37] (for the low-frequency tail). The study further uncovers a new potential complementarity between GW detection and laboratory experiments at energy, intensity, and neutrino frontiers, such as the LHC, DUNE, SHiP, and FASER [38–45], in search for new physics, depending on the details of the asymmetry transfer interactions between the AD field and the SM.

Benchmark models.

AD scalars: Background dynamics and asymmetry generation: We consider a complex scalar field Φ (SM gauge singlet) with canonical kinetic terms, Einstein gravity, and a simple renormalizable polynomial potential,

$$V(\Phi) = \lambda_\Phi |\Phi|^4 + m_\Phi^2 |\Phi|^2 - A(\Phi^n + \Phi^{*n}), \quad (1)$$

where $n \leq 4$ is an integer. We choose $n = 2$ for concreteness. Although other choices may lead to similar results, $n = 2$ stands out as a simple and natural option [46]. The first two terms manifest a $U(1)$ symmetry, which is broken by the third term as necessary to generate the $\Phi\text{--}\Phi^\dagger$ asymmetry, a precursor to the baryon asymmetry. We consider the general possibility that Φ , as a light spectator during inflation, starts with a large (even Planckian) field displacement [48–52], for instance, due to de Sitter fluctuations during inflation. We have verified that the isocurvature constraints from cosmic microwave background (CMB) data, which are relevant for any light spectator field, can be satisfied (see Sec. B of Supplemental Material for details [53]). We denote the typical value at the end of inflation as Φ_{in} , which is the initial value for our calculations [82]. As long as the (radial) field is light, $H \gg H_{\text{osc}} \simeq \sqrt{V_{,|\Phi|}/|\Phi|}$, where H_{osc} is the effective mass of $|\Phi|$ [49], it remains frozen [83]. The initial phase of Φ , $\theta = \text{Arg}(\Phi)$, is also set as θ_{in} at the end of inflation in our observable patch of the Universe. As the Universe expands, the Hubble rate drops, and at $H \simeq H_{\text{osc}}$ the field starts rolling. For our parameters, the rolling starts during the radiation-dominated (RD) era, well after the end of inflation.

The real and imaginary parts of $\Phi \equiv \frac{1}{\sqrt{2}}(\phi_R + i\phi_I)$ obey the following equations of motion:

$$\ddot{\phi}_R + 3H\dot{\phi}_R + m_R^2\phi_R + \lambda_\Phi(\phi_R^2 + \phi_I^2)\phi_R = 0, \quad (2)$$

$$\ddot{\phi}_I + 3H\dot{\phi}_I + m_I^2\phi_I + \lambda_\Phi(\phi_R^2 + \phi_I^2)\phi_I = 0, \quad (3)$$

where $m_{R,I}^2 = m_\Phi^2 \mp 2A$. The Hubble expansion rate in the RD Universe is $H \propto t^{1/2}$. An overdot denotes derivative with respect to cosmic time t . If $A \neq 0$, the real and imaginary components have different masses, so their trajectories in the $\phi_R - \phi_I$ plane resemble shrinking ellipses (not straight lines), even for zero initial velocities.

The background field evolution (assuming $A \ll m_\Phi^2$) is initially dominated by the $|\Phi|^4$ term, where the amplitude scales as $\phi \propto 1/a$. When the amplitude drops below $|\Phi_*| = m_\Phi/\sqrt{\lambda_\Phi}$ at $t = t_*$, the $|\Phi|^2$ term dominates the oscillation and the fields evolve approximately as $\phi_{R,I} \sim a^{-3/2} \cos(m_{R,I}t + c_{R,I})$, with $c_{R,I}$ being a constant phase (see Supplemental Material, Sec. C, for details). The $\Phi - \Phi^\dagger$ asymmetry, which will be subsequently transferred to the baryon (or lepton) number, is $n_\Phi(t) \equiv i(\Phi^\dagger \dot{\Phi} - \dot{\Phi}^\dagger \Phi) = \dot{\phi}_R \phi_I - \phi_R \dot{\phi}_I$. In the absence of significant washout, n_Φ directly corresponds to n_B or n_L (see the Discussion later). Using the above approximations for $\phi_{R,I}$ and including the terms $\Gamma_\Phi \dot{\phi}_{R,I}$ to model the decay of Φ to SM states (the specific form of Γ_Φ depends on the decay channels, which we elaborate on later), the comoving baryon asymmetry generated during the time $t > t_*$ is [47,84–86]

$$\begin{aligned} \left(\frac{a(t)}{a_{\text{in}}}\right)^3 n_B(t) &\simeq 4A\phi_{R,\text{in}}\phi_{I,\text{in}} \left(\frac{\phi_{\text{in}}}{\phi_*}\right) \\ &\times \int_{t_*}^t dt' \cos(m_R(t' - t_*)) \sin(m_I(t' - t_*)) \\ &\times e^{-\Gamma_\Phi(t' - t_*)}, \end{aligned} \quad (4)$$

which can be further simplified for $A \ll m_\Phi^2$ (the detailed derivation is presented as Supplemental Material) and scales as

$$\frac{a^3(t)}{a_{\text{in}}^3} n_B(t) \propto \int_{t_*}^t dt' \sin\left(\frac{2A}{m_\Phi}(t - t_*)\right) e^{-\Gamma_\Phi(t - t_*)}. \quad (5)$$

Evaluating the above integral for $t \rightarrow \infty$ and factoring in the expansion during the RD Universe, we express n_B/s as

$$\frac{n_B}{s} = \begin{cases} \left(\frac{4\alpha^3}{\lambda_\Phi k_{T_d}^5}\right)^{1/4} \epsilon_\Phi \frac{m_\Phi M_{\text{Pl}}}{T_d^2} \sin(2\theta); & \gamma_\Phi \gg 2\epsilon_\Phi, \\ \left(\frac{\alpha^3 k_{T_d}^2}{64\lambda_\Phi}\right)^{1/4} \frac{1}{\epsilon_\Phi} \frac{T_d^2}{m_\Phi M_{\text{Pl}}} \sin(2\theta); & \gamma_\Phi \ll 2\epsilon_\Phi, \end{cases} \quad (6)$$

where $M_{\text{Pl}} \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass, and we define the dimensionless quantities $\epsilon_\Phi = A/m_\Phi^2$, $\gamma_\Phi = \Gamma_\Phi/m_\Phi$, and $k_{T_d} = \sqrt{\pi^2 g(T_d)/90}$. T_d is the temperature of the Universe characterizing Φ decay, as defined by $H(T_d) = \Gamma_\Phi$, α is the initial fraction of energy density in the spectator sector Φ relative to the total energy density, and $\theta = \theta_{\text{in}}$ is defined earlier. The exact form of n_B/s , from which Eq. (6) is derived, is given as Supplemental Material.

The change in the scaling of the final asymmetry with the normalized symmetry-breaking parameter $\epsilon_\Phi \equiv A/m_\Phi^2$ shown by the two branches of Eq. (6) can be understood simply in terms of the competition between two timescales: the oscillation period of the asymmetry and the lifetime of the scalar. The symmetry-breaking parameter ϵ_Φ induces a mass splitting between the particle and antiparticle components of Φ , leading to an oscillation of the charge

asymmetry with frequency $\omega_\Phi \sim \Delta m_\Phi \sim \epsilon_\Phi m_\Phi$ and period $T_{\text{asy}} \sim 1/(\epsilon_\Phi m_\Phi)$. At the same time, the scalar decays with rate Γ_Φ , corresponding to a lifetime $\tau_\Phi \sim 1/\Gamma_\Phi$. The net asymmetry transferred to the thermal bath is proportional to the time integral of an oscillatory source with an exponential envelope, schematically $n_B \propto \int_0^\infty dt e^{-\Gamma_\Phi t} \sin(\omega_\Phi t)$, so the final result depends on whether the decay occurs before or after many oscillations. In the regime $\Gamma_\Phi \gg \omega_\Phi$ (i.e., $\tau_\Phi \ll T_{\text{asy}}$), the integral is dominated by very early times, when the sine function has not yet completed a full oscillation and may be approximated as $\sin(\omega_\Phi t) \simeq \omega_\Phi t$. The generated asymmetry is then essentially fixed before sign-flipping oscillations set in, and scales linearly with the oscillation frequency (equivalently with the asymmetry parameter ϵ_Φ) as

$$n_B \propto \int_0^{\sim 1/\Gamma_\Phi} dt e^{-\Gamma_\Phi t} \omega_\Phi t \sim \frac{\omega_\Phi}{\Gamma_\Phi} \propto \epsilon_\Phi. \quad (7)$$

By contrast, in the regime $\Gamma_\Phi \ll \omega_\Phi$ (i.e., $\tau_\Phi \gg T_{\text{asy}}$), the source oscillates many times before the field decays, spending comparable time with positive and negative signs. As a result, contributions from successive oscillations partially cancel, and the net asymmetry is a residue arising from the slow decay of the envelope. To see the overall scale of the asymmetry for this limit, we note that the contribution from a single half-period of the oscillation is of order $\int dt \sin(\epsilon_\Phi m_\Phi t) \sim 1/(\epsilon_\Phi m_\Phi)$, which sets the characteristic scale of the integral in Eq. (5). Including all subsequent oscillations leads to an alternating geometric series, which can be summed and consequently multiplies the result of the first oscillation by $1/2$ for the case with $\gamma_\Phi/2\epsilon_\Phi \ll 1$. The final asymmetry thus scales as

$$n_B \propto \int_0^\infty dt \sin(\epsilon_\Phi m_\Phi t) \simeq \frac{1}{2} \times \int_0^{\pi/(2\epsilon_\Phi m_\Phi)} dt \sin(\epsilon_\Phi m_\Phi t) \sim \frac{1}{2\epsilon_\Phi m_\Phi}. \quad (8)$$

In summary, a quadratic symmetry-breaking term induces a constant mass splitting between the real and imaginary components of Φ , leading to a single well-defined oscillation frequency $\omega_\Phi \propto \epsilon_\Phi m_\Phi$. This simple harmonic structure makes the final asymmetry an explicit function of ω_Φ , and hence of ϵ_Φ , producing the characteristic two-regime behavior [87] $n_B \propto \epsilon_\Phi$ for $\Gamma_\Phi \gg \omega_\Phi$ and $n_B \propto 1/\epsilon_\Phi$ for $\Gamma_\Phi \ll \omega_\Phi$, shown in Eq. (6). In the case where the asymmetry in the Φ field is first transferred to an SM lepton asymmetry n_L , which is later converted to a baryon asymmetry n_B via the sphaleron process, the above result gets multiplied by a factor of $28/79$ since [88–90] $n_B/s = -(28/79) \times n_L/s$.

Equation (6) is derived assuming that the asymmetry is generated when the potential is dominated by the $|\Phi|^2$ term, thus requiring that the background dynamics prior to t_* is not affected by the A term. The corresponding condition on A is derived as follows: the only term in the potential that

involves the angular degree of freedom is the symmetry-breaking term $2A|\Phi(t)|^2 \cos 2\theta(t) \subset V(|\Phi|, \theta)$, indicating an upper bound on the effective mass along the angular direction: $|m_\theta^2| \leq 2A$. The angular degree of freedom is frozen at θ_{in} for $t < t_*$ if $|m_\theta^2| < H^2(t_*)$ [85], which translates into $\epsilon_\Phi \equiv A/m_\Phi^2 < (\mu^2/16\alpha)(\phi_{\text{in}}^2/M_{\text{Pl}}^2)$, where $\mu^2 \equiv m_\Phi^2/(\lambda_\Phi \phi_{\text{in}}^2)$, and α is defined below Eq. (6). As we see in the next section, $\phi_{\text{in}} \simeq M_{\text{Pl}}$ and $\mu^2 \lesssim 10^{-5}$ are required for detectable GW signals. With larger A , motion along the angular direction is non-negligible, and $\theta(t)$ would have already evolved toward its minimum before $t = t_*$, thus reducing $\sin \theta$ and suppressing the asymmetry [see Eq. (6)]. For our benchmark examples, the chosen A values satisfy this self-consistency condition.

We now present some numerical examples of the asymmetry. For $\epsilon_\Phi = 10^{-20}$ the observed $n_B/s \sim 10^{-10}$ can be realized with $T_d \sim 2.5 \times 10^7$ GeV when $m_\Phi = 10$ GeV. As we will discuss, transferring the asymmetry to the baryon sector may involve a significant (model-dependent) washout, requiring a larger initial Φ asymmetry to compensate. Therefore, we also consider parameter points that lead to $n_\Phi/s \gg 10^{-10}$, for example, the same ϵ_Φ with $T_d \sim 1$ GeV gives $n_\Phi/s \sim 4 \times 10^{-5}$. A more detailed parameter scan is provided in Sec. C of Supplemental Material, demonstrating the potential to produce a large initial asymmetry that compensates for possible washout effects.

Transfer to the SM baryon asymmetry: The asymmetry in the Φ field generated during its oscillation needs to eventually be transferred to the SM baryon asymmetry via its subsequent decay to baryons or leptons (further assisted by electroweak sphalerons in the latter case). This can generally be realized by a B- or L-violating coupling involving Φ in terms of effective higher-dimensional operators, such as $\Phi Q Q Q L/\Lambda^3$ [6], $\Phi U D D U D D/\Lambda^6$, and $\Phi L L L L L E/\Lambda^6$. Whether the related constraints, for instance, from $n - \bar{n}$ oscillation, can be satisfied depends on the specifics of the UV completion, including flavor structure in the couplings. As a proof of principle, we outline examples of renormalizable, UV-complete models that realize the Φ to baryon asymmetry transfer in the following, and summarize the conditions on viable model parameters. Further details are given in Sec. D of Supplemental Material. The core results of this work are independent of the details of the asymmetry transfer mechanism.

(A) *Φ decay to leptons (leptogenesis).*

A simple renormalizable model involves

$$\mathcal{L} \supset y_{N,i} \Phi \bar{N}_i^C N_i + g_{\nu,i} \bar{N}_i H L_j + \text{H.c.}, \quad (9)$$

where N_i 's are right-handed Majorana neutrinos (with $i = 1, 2, 3$ as a typical choice), H is SM Higgs, and L_j are SM lepton doublets. This model potentially relates to neutrino mass generation, which, nevertheless, is not imposed as a requirement here. We consider the simple scenario where Φ

decays to NN , then N transfers the asymmetry to SM leptons via freeze-out or freeze-in processes enabled by the NHL coupling, as considered in the literature of leptogenesis [91–97].

(B) Φ decay to baryons (direct baryogenesis).

A simple renormalizable model, in this case, includes

$$\mathcal{L} \supset y_\chi \Phi \chi \chi + g_i \chi \bar{u}_i \varphi + \lambda_{ij} \varphi d_i d_j + \text{H.c.}, \quad (10)$$

where χ is a singlet Majorana fermion, φ is an up-type diquark scalar, and the flavor indices $i, j = 1, 2, 3$ in down-type quark combination $d_i d_j$ should be antisymmetric. In the viable parameter region, the decay proceeds with $\Phi \rightarrow \chi \chi$, followed by $\chi \rightarrow udd$.

Both models are subject to cosmological conditions that impose lower limits on m_Φ , m_χ , or m_N . For Φ to decay to baryons [before big bang nucleosynthesis (BBN)], it is subject to kinematic conditions: $m_\Phi \gtrsim 2 \text{ GeV}$, and $m_\chi \gtrsim 1 \text{ GeV}$. For Φ asymmetry to be transferred to SM leptons via N before the electroweak (EW) phase transition requires $m_\Phi > 2m_N \gtrsim O(0.1) \text{ GeV}$ [92]. In addition, for both models, the potential asymmetry washout effects should be either suppressed or compensated by a larger initial Φ asymmetry produced during its oscillation (which has been shown to be viable in the last section). Furthermore, relevant laboratory experiments at the energy and intensity frontiers, searching for B-violation, color-charged new particles, and sterile neutrinos, among others, may provide additional constraints on model parameters. In Sec. D of Supplemental Material, we elaborate on relevant constraints and demonstrate examples of viable parameter regions.

Parametric resonance and GW signals. We have seen that the generation of the baryon asymmetry is determined by the classical motion of Φ as an averaged background field. In addition to this classical motion of Φ , the corresponding fluctuations around it can lead to parametric resonance and source GWs. We decompose $\phi_{R,I} = \bar{\phi}_{R,I}(t) + \delta\phi_{R,I}(x, t)$, where $\bar{\phi}_{R,I}$ describe the classical background motion and $\delta\phi_{R,I}(x, t)$ correspond to the fluctuations, which have a quantum origin. The evolution of background fields $\bar{\phi}_{R,I}$ follows Eqs. (2) and (3), whereas we expand the fluctuations in Fourier modes as

$$\delta\phi_R(x, t) = \int \frac{d^3k}{(2\pi)^3} [a_{R,k} e^{i\vec{k}\cdot\vec{x}} \delta\phi_{R,k}(t) + \text{H.c.}] \quad (11)$$

and similarly for $\delta\phi_I(x, t)$.

As the background fields oscillate around the minimum of the potential, the fluctuations in momentum space evolve according to

$$\begin{aligned} \ddot{\delta\phi}_{R,I} + 3H\dot{\delta\phi}_{R,I} + m_{R,I}^2 \delta\phi_{R,I} + \frac{k^2}{a^2} \delta\phi_{R,I} \\ + \lambda_\phi (3\bar{\phi}_{R,I}^2 + \bar{\phi}_{I,R}^2) \delta\phi_{R,I} + 2\bar{\phi}_I \bar{\phi}_R \delta\phi_{I,R} = 0. \end{aligned} \quad (12)$$

We thus see that, in general, the equations for $\delta\phi_{R,I}$ are coupled. We first consider the representative limit of $A = 0$, where $m_R = m_I = m_\Phi$, and the background field trajectory is simply a straight line through the origin (for vanishing initial velocities, as is the case here).

In this case, we can rotate our field basis and define fields in the radial direction (along the background motion) $\phi_{\parallel}$ and perpendicular to it (orthogonal direction) as $\phi_{\perp}$. By construction, $\bar{\phi}_{\perp} = 0$, and the fluctuation equations for $\phi_{\parallel}$ and $\phi_{\perp}$ decouple, namely,

$$\begin{aligned} \ddot{\delta\phi}_{\parallel,\perp} + 3H\dot{\delta\phi}_{\parallel,\perp} + \frac{k^2}{a^2} \delta\phi_{\parallel,\perp} + m_\Phi^2 \delta\phi_{\parallel,\perp} \\ + q_{\parallel,\perp} \lambda_\Phi \bar{\phi}_{\parallel,\perp}^2 \delta\phi_{\parallel,\perp} = 0, \end{aligned} \quad (13)$$

where $q_{\parallel,\perp} = \{3, 1\}$. Quasiperiodic background field oscillations can exponentially amplify fluctuations within certain wave number ranges. This rapid growth of mode occupation numbers allows these modes to be treated classically [98,99]. This phenomenon, originally studied for exponential particle production (“preheating”) after inflation [100], can also occur for any oscillating scalar (see, for instance, [49]).

To quantify parametric resonance, we define a dimensionless field amplitude $\tilde{\phi} = \phi/\phi_{\text{in}}$ and a corresponding conformal scaling for time as $d\eta = \sqrt{\lambda_\Phi} \phi_{\text{in}} dt/a$, with ϕ_{in} being the initial amplitude of the scalar field $\phi_{\text{in}} = |\Phi_{\text{in}}|$. The dimensionless fluctuation equation becomes ($\delta\phi' \equiv d\delta\phi/d\eta$)

$$\delta\phi''_{\parallel,\perp} + 2\tilde{H}\delta\phi'_{\parallel,\perp} + (\tilde{k}^2 + a^2\mu^2 + q_{\parallel,\perp} a^2 \tilde{\phi}^2) \delta\phi_{\parallel,\perp} = 0, \quad (14)$$

where $\tilde{k}^2 = k^2/\lambda_\Phi \phi_{\text{in}}^2$ and $\tilde{\phi} = \phi/\phi_{\text{in}}$.

There are two essential resonance parameters: $\mu^2 = \frac{m_\Phi^2}{\lambda_\Phi \phi_{\text{in}}^2}$ and $q_{\parallel,\perp}$. As shown in [100], a large μ^2 can suppress the resonance and determine if the system displays *stable resonance*, *stochastic resonance*, or *no decay*. We choose to focus on $\mu \leq 10^{-5}$, such that the system is always in a stable resonance regime. For such small μ^2 , the system is similar to the well-studied case of a quartic field during preheating, whereas $q_{\parallel,\perp}$ is fixed in our case due to the rotational symmetry of the potential for $A = 0$. Section E of Supplemental Material discusses in detail the effect of $q_{\parallel,\perp}$ on the growth of $\delta\phi_{\parallel,\perp}$ and shows that for sufficiently small values of the asymmetry parameter $A \lesssim 0.005 m_\Phi^2$, the parametric resonance dynamics is indistinguishable from the case $A = 0$, and thus the GW amplitude remains insensitive to the value of A . Fortunately, the parameter range of A that produces sufficient baryon asymmetry, as considered in our benchmarks, is well within this regime. On a technical note,

we use the parametrization of real and imaginary components of Φ (ϕ_R , ϕ_I) when computing the baryon asymmetry, whereas we are using the radial direction and its orthogonal ($\phi_{\parallel}$, $\phi_{\perp}$) when computing parametric resonance for $A \ll m_{\Phi}^2$. This is due to the fact that the baryon asymmetry scales with A and thus the difference between the ϕ_I and ϕ_R directions is important, whereas parametric resonance and GW production are blind to $A \ll m_{\Phi}^2$, making radial coordinates more suitable. Interesting deviations from the $A = 0$ limit are discussed in Sec. E of Supplemental Material. For $A \gtrsim 0.05 m_{\Phi}^2$, the GW spectrum starts being sensitive to the exact value of A as well as the initial angle θ_{in} . This provides intriguing possibilities for distinguishing between these parameters by dedicated analysis of the GW spectral shape.

The growth of fluctuations and their backreaction breaks the homogeneous oscillation of Φ , leading to inhomogeneities with an effective anisotropic stress tensor $\Pi_{ij}^{TT} = \nabla_i \phi^a \nabla_j \phi^a$ (where the repeated index “ a ” indicates the number of the scalar components and a summation is implied). The dynamics of the generated GWs are governed by the linearized equation for the transverse-traceless tensor perturbation sourced by the above anisotropic stress,

$$\ddot{h}_{ij} + 3H\dot{h}_{ij} - \frac{\nabla^2 h_{ij}}{a^2} = \frac{2}{a^2 M_{\text{Pl}}^2} \Pi_{ij}^{TT}. \quad (15)$$

The quantity of interest is the energy density power spectrum of the GWs normalized by the critical energy density of the Universe today. For GWs produced in a radiation-dominated era, the spectrum of the energy density of GWs (per logarithmic momentum interval) observable today is [13,49,101–105]

$$\begin{aligned} \Omega_{\text{GW},0} h^2 &= \frac{h^2}{\rho_{\text{crit}}} \frac{d\rho_{\text{GW}}}{d \ln k} \Big|_{t=t_0} = \frac{h^2}{\rho_{\text{crit}}} \frac{d\rho_{\text{GW}}}{d \ln k} \Big|_{t=t_e} \frac{a_e^4 \rho_e}{a_0^4 \rho_{\text{crit},0}} \\ &= \Omega_{\text{rad},0} h^2 \Omega_{\text{GW},e} \left(\frac{g_*}{g_0} \right)^{-1/3}, \end{aligned} \quad (16)$$

where the critical density of the Universe is $\rho_{\text{crit}} = 3M_{\text{Pl}}^2 H^2$ and $\Omega_{\text{rad},0} h^2 = h^2 \rho_{\text{rad},0} / \rho_{\text{crit}} \simeq 4.3 \times 10^{-5}$. In particular, we employ the publicly available code Cosmolattice [106,107] in a three-dimensional lattice, in order to numerically solve the metric perturbation equations in Eq. (15) along with the equations for the scalar system in an expanding Universe.

The spectra for three benchmark examples are shown in Fig. 1. For a Φ^4 -dominated system, as in our case initially, the field starts rolling/oscillating when $H \sim \sqrt{\lambda_{\Phi}} \phi_{\text{in}}$. The combination $\sqrt{\lambda_{\Phi}} \phi_{\text{in}}$ also determines the frequency of the observed GWs. We consider the RD universe with the scalar field system contributing a fraction of the total energy α . Figure 1 corresponds to $\alpha = 0.3$. For GWs generated by sources peaked around a characteristic scale k_c in momentum space [108], $\Omega_{\text{GW}}^{\text{peak}} \propto \alpha^2 \propto \phi_{\text{in}}^8$, and we found that a sizable $\alpha \sim O(0.1)$, which optimizes GW amplitude, requires $\phi_{\text{in}} \sim M_{\text{Pl}}$ (see also [49]). Furthermore, the observed frequency

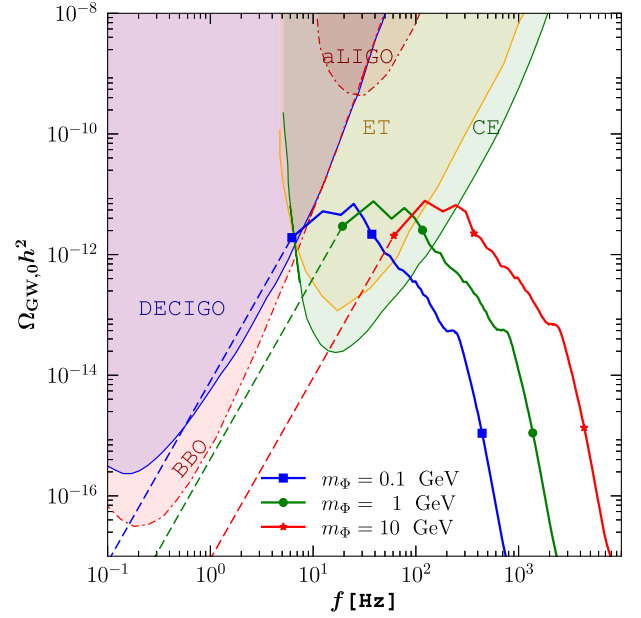


FIG. 1. We show the GW spectrum originating from the AD model considered here for three benchmark masses of Φ . The potential parameters $\lambda_{\Phi} \sim 10^{10} m_{\Phi}^2 / M_{\text{Pl}}^2$, $A = 0$ are chosen to ensure efficient parametric resonance (see the main text for details). The fraction of initial scalar field energy density over the total energy density of the Universe is $\alpha = 30\%$. The dashed lines at lower frequencies are extrapolated according to the causal superhorizon k^3 scaling.

today is $f_{\text{GW}}^{\text{peak}} \simeq 2.7 \times 10^{10} k_c / \sqrt{H_c M_{\text{Pl}}} \text{ Hz}$, where $k_c \sim H \sim \sqrt{\lambda_{\Phi}} \phi_{\text{in}}$ (see Sec. E of Supplemental Material for details), leading to $f_{\text{GW}}^{\text{peak}} \sim 10^{10} \lambda_{\Phi}^{1/4}$ for $\phi_{\text{in}} \sim M_{\text{Pl}}$. As we have seen from the discussion on the resonance structure of the model, the parameters $\mu_{\pm}$ (or simply μ when $A = 0$) determine the nature and efficiency of resonance [100]. We found that stable resonance occurs for $\mu \equiv m_{\Phi} / (\sqrt{\lambda_{\Phi}} \phi_{\text{in}}) \lesssim 10^{-5}$. We choose $\mu = 10^{-5}$, which fixes the product $\lambda_{\Phi} \phi_{\text{in}}^2$. With an optimal initial value of $\phi_{\text{in}} \sim M_{\text{Pl}}$, we find that $\lambda_{\Phi} \simeq 10^{-35}$, 10^{-33} , and 10^{-31} for $m_{\Phi} = 0.1$, 1, and 10 GeV, see Fig. 1. As discussed earlier, $m_{\Phi} \gtrsim O(0.1) \text{ GeV}$ is required for Φ to successfully transfer the asymmetry to SM baryons. Interestingly, in the viable range of $m_{\Phi} \sim O(0.1) - O(10) \text{ GeV}$, the resulting peak GW frequency lies well within the range of the ET/CE and with sufficient amplitudes, making signal detection promising for the near future. A higher range of m_{Φ} is compatible with baryogenesis but requires futuristic, higher-frequency GW detectors [109] to capture the peak frequency range of the signal. Due to the finite box size and lattice spacing, low-frequency modes well off the peak cannot be resolved in our simulation, which is a common issue in the lattice simulation literature [13]. Therefore, to extend the spectral prediction to lower frequencies, we apply the standard assumption that at low wave numbers, the GW spectrum scales as k^3 , as derived from causality arguments [13,110–112]. This is shown in dashed

lines in Fig. 1. Depending on m_Φ , the low- k tail may be detectable by DECIGO and BBO, providing the prospect for correlated detection across multiple frequency bands, which simultaneously probes the peak and tail of the spectrum. Before concluding, let us note that the simulated signal shown in Fig. 1 is $\mathcal{O}(100)$ times larger than the lowest point of the CE sensitivity curve. Since $\Omega_{\text{GW}} h^2 \propto \alpha^2$ [108], we can reduce α , the initial fraction of the energy density of the scalar field with respect to the total energy density in the Universe, down to $\alpha \sim \mathcal{O}(1\%)$ and still yield a GW signal within reach of the CE.

Conclusion and outlook. In this Letter, we first demonstrate that the Affleck-Dine mechanism for baryogenesis can be realized by a complex light scalar field Φ with a mass m_Φ well below the electroweak scale without invoking supersymmetry. In particular, we use a convex, renormalizable, polynomial potential (which does not lead to the formation of Q-balls, generally present in SUSY AD models). Such a scalar field acts as a spectator field during inflation while acquiring a large field value at the end of inflation. The subsequent rolling down and oscillation of the scalar field occur long after inflation, during the radiation-dominated era, triggering the generation of baryon asymmetry. Parametric resonance and fragmentation generally arise during the oscillation of the scalar field, naturally leading to GW production. We demonstrate that the models considered can lead to successful baryogenesis while sourcing detectable GW signals with a peak frequency of $\mathcal{O}(10\text{--}100)$ Hz, within reach of experiments such as the ET and CE. Furthermore, the characteristic new physics scale, characterized by m_Φ , intriguingly lies in the range of $\mathcal{O}(0.1\text{--}10)$ GeV, while

the transfer of the Φ asymmetry to the SM B- or L-asymmetry requires interactions between the new physics sector and the SM states. Hence, from this well-motivated scenario, a new, natural complementarity arises between SGWB detection and laboratory searches for new particle physics across the energy, intensity, and neutrino frontiers. The specifics of the complementary laboratory signal depend on the details of the asymmetry transfer mechanism. Examples include searches for $\mathcal{O}(0.1\text{--}10)$ GeV mass sterile neutrino with DUNE, SHiP, FASER, etc., and exotic multijet signals at the LHC (see Supplemental Material, Sec. D). This work thus sheds new light on probing early Universe baryogenesis through its GW imprint, potentially complemented by associated laboratory probes, while presenting a well-motivated new physics source for SGWB.

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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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