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Primordial Black Holes from Inflaton Fragmentation into Oscillons

Eric Cotner,[1,*] Alexander Kusenko,[1,2,†] and Volodymyr Takhistov[1,‡]
[1] *Department of Physics and Astronomy, University of California, Los Angeles Los Angeles, CA 90095-1547, USA*
[2] *Kavli Institute for the Physics and Mathematics of the Universe (WPI), UTIAS The University of Tokyo, Kashiwa, Chiba 277-8583, Japan*
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We show that fragmentation of the inflaton into long-lived spatially localized oscillon configurations can lead to copious production of black holes. In a single-field inflation model primordial black holes of *sublunar* mass can form, and they can account for all of the dark matter. We also explore the possibility that solar-mass primordial black holes, particularly relevant for gravitational wave astronomy, are produced from the same mechanism.

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Primordial black holes (PBHs) can form in the early Universe and can account for all or part of the dark matter (DM) (e.g. [1–17]). They have also been linked to a variety of topics in astronomy, including the recently discovered [18–20] gravitational waves [21–28], formation of supermassive black holes [22, 29, 30] as well as r-process nucleosynthesis [31] and gamma-ray bursts [32] from compact star disruptions.

Many proposed scenarios of PBH formation assume that inflation has generated some excess of density perturbations on certain scales, which produce PBHs when they re-enter the Hubble horizon during the radiation dominated phase or during some intermediate matter-dominated stage (for review, see [5, 33]). The required inflaton potentials could be *ad hoc*, or can be well-motivated in the context of hybrid inflation [4], supergravity [7], etc. PBHs can also form from large extended objects, such as non-topological solitons in supersymmetric theories, which behave as matter and come to dominate the universe for a short time before decaying [8, 34]. In this case, the overdensities needed for PBH formation result from statistical fluctuations in a system with a relatively low number of very massive "particles" and not from the spectrum of primordial density fluctuations. In this *Letter* we show that if the inflaton potential admits long-lived oscillon solutions, their formation can lead to copious production of PBHs.

Oscillons [35–40] arise in many well motivated theories with scalar fields, such as models of inflation [41], axions [42] or moduli [43]. The oscillons are localized, metastable, pseudo-solitonic configurations of real scalar fields. The stability of an oscillon is not guaranteed by a conserved charge and its long lifetime is associated with an approximate adiabatic invariant [39, 44]. Early Universe oscillons have been recently studied in connection with primordial gravity waves [45] as well as baryogenesis [46].

For definiteness, we consider the model with a single inflaton field ϕ that has a canonical kinetic term, minimal coupling to Einstein gravity and a potential [41, 47–49]

$$V = \frac{m^2}{2}\phi^2 - \frac{\lambda}{4}\phi^4 + \frac{g^2}{6m^2}\phi^6 \ . \tag{1}$$

Here $\lambda > 0$ and for convenience, following the original studies, we take $\lambda, g, m/M_{\rm pl} \ll 1$, where $M_{\rm pl}$ is the Planck mass. The model [41, 47–49] is inspired by a class of well motivated supergravity and string theories [41, 50, 51] and could be considered as a Taylor series expansion of a more general potential for some range of the scalar field. We assume that the density perturbations that seed cosmological structures are generated when the inflaton field has a much larger value, for which the shape of the potential is not necessarily described by Eq. (1). After the inflationary phase, the inflaton begins to oscillate near the minimum of the potential as described by Eq. (1). At this time, the inflaton condensate fragments into oscillons.

A necessary condition for oscillon formation is that the potential is shallower than quadratic near the minimum (making the scalar self-interactions attractive). This is the case for $\lambda > 0$. For $(\lambda/g)^2 \ll 1$, the above potential admits "flat-top" oscillons, which are extremely stable on the cosmological time scales and for which the analytic description is known [41, 47–49].

An initially homogeneous inflaton condensate can fragment into lumps, corresponding to oscillons. The inflaton self-resonance parametrically amplifies field fluctuations $\delta\phi_k$ in some band of wave-numbers k around the background field $\overline{\phi}$. This can be analytically investigated through Floquet analysis, where the most unstable modes behave as $\delta\phi_k(t) \propto e^{\mu_k t}P(t)$, with μ_k denoting the Floquet exponent and $P(t)$ a periodic function. In an expanding background significant amplification of fluctuations requires $\mu_k(a)/H \gg 1$, where $a(t)$ is the cosmic scale factor and $H = H_i/\sqrt{a^3}$ is the Hubble parameter at the bottom of potential. At fragmentation $H_i \simeq \sqrt{\lambda/10g^2}(m/M_{\rm pl})m$ and $a_i = 1$. The amplification condition translates to [47]

$$\frac{\mu_k(a)}{H} = \frac{M_{\rm pl}}{m}\Big(\frac{\lambda^{3/2}}{g}\Big)\Big[\sqrt{\frac{9}{4}\frac{\tilde{k}^2}{a^2}\Big(1-\frac{1}{a^3}\Big)-\Big(\frac{\tilde{k}^4}{a}\Big)}\Big] \gg 1\ , \tag{2}$$

where $\tilde{k} = (g/\lambda m)k$, with k being related to the physical wavenumber via $k = a\,k_p$. The total amplification of fluctuations as they pass through the instability band is found by integrating over the Floquet exponent as

$$\delta\phi_k(a) \sim \frac{1}{\sqrt{2\omega_k}}\frac{1}{a^{3/2}}e^{\beta f(\tilde{k},a)}\ , \tag{3}$$

where

$$f(\tilde{k},a) = \sqrt{\frac{5}{2}}\int_C d\log\overline{a}\Big[\tilde{k}\sqrt{\frac{9}{10\overline{a}^2}\Big(1-\frac{1}{\overline{a}^3}\Big)-\frac{\tilde{k}^2}{\overline{a}}}\Big] \tag{4}$$

and $C = \frac{a^3-1}{a} > \frac{10}{9}\tilde{k}^2$, $\beta = \sqrt{\lambda}(\lambda/g)(M_{\rm pl}/m)$ and $\omega_k^2 \simeq k^2+m^2$. Since $\beta \sim \mu/H$, we are interested in the $\beta \gg 1$ regime.

The condition [47] for formation of oscillons from amplified perturbations can be formulated as $k^{3/2}\delta\phi \sim \overline{\phi}$. The average number density of oscillons can then be estimated as $\overline{n} \sim (k_{nl}/2\pi)^3/a^3$, where $k_{nl} \sim \beta^{-1/5}(\lambda/g)m$ label the modes that become non-linear the earliest. While the model supports several distinct oscillon populations, we focus on the flat-top oscillons due to their stability. Unlike the usual Gaussian-profile oscillons, they possess an approximately uniform core density of $\rho_c \simeq m^4(9\lambda/20g^2)$ [48]. Taking the characteristic radius of oscillons to be $R \sim \pi/k_{nl}$, we can estimate their energy as

$$E \sim \frac{4\pi\rho_c R^3}{3} \simeq \frac{3\pi^4 m^4}{5k_{nl}^3}\Big(\frac{\lambda}{g^2}\Big)\ . \tag{5}$$

The above provides $\overline{n}(E)$ through k_{nl} substitution.

Given an average number density $\overline{n}$ of uniformly distributed objects, the probability of finding N objects within a volume V follows the Poisson distribution

$$P_N(N) = \frac{(\overline{n}V)^N}{N!}e^{-\overline{n}V}\ . \tag{6}$$

The total mass of a cluster of oscillons is $M = NE$. Hence, the probability distribution of oscillon cluster mass is given by $P_M(M) = \sum_N P_N(N)\delta(M-NE)$. The delta-function can be eliminated through a Fourier transform

$$\tilde{P}_M(\mu) = \int dM\, P_M(M)\, e^{iM\mu} = e^{\overline{n}V(e^{iE\mu}-1)}\ , \tag{7}$$

followed by an inverse transform

$$P_M(M) = \frac{1}{2\pi}\int d\mu\, e^{-iM\mu}\, e^{\overline{n}V(e^{iE\mu}-1)}\ . \tag{8}$$

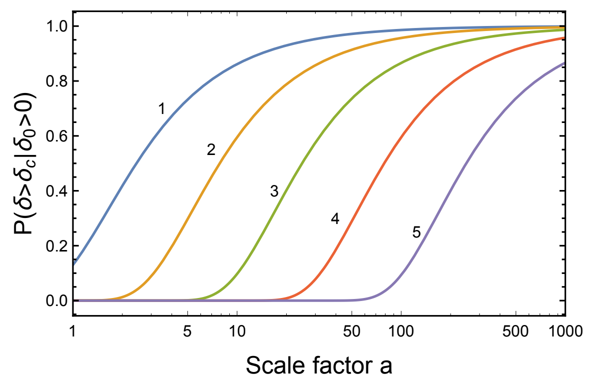

FIG. 1: Fraction of super-critical overdensities as a function of the scale factor $a(t)$. Results are shown for several different volumes, which are expressed in terms of the Hubble horizon as $V/V_H = 3\times10^{-5}$, 3×10^{-4}, 3×10^{-3} , 3×10^{-2} , 3×10^{-1} and labeled “1”, “2”, “3”, “4” and “5”, respectively. The values of the input parameters $(\lambda/g)^2 = 0.2$ and $\beta = 56.1$ correspond to those of Model A.

An approximate analytic non-integral form of Eq. (8) can be found through the method of steepest descent. The resulting expression is

$$P_\eta(\eta) = \frac{1}{\sqrt{2\pi\beta^{3/5}\eta}}e^{\beta^{-3/5}[\eta(1-\ln\{(2\pi)^3\eta/v\})-v/(2\pi)^3]}\ , \tag{9}$$

where $v = V(\lambda m/g)^3 = V[M_{\rm pl}(\lambda/g)(m/M_{\rm pl})]^3$ and $\eta = M/E_0$, with $E_0 = (3\pi^4/5)[M_{\rm pl}(\lambda/g)/\beta^2(m/M_{\rm pl})]$, denote the rescaled dimensionless volume and mass, respectively.

Using $P_\eta(\eta)$ we can now calculate the distribution of the initial density contrasts δ_0. In terms of η, δ_0 is

$$\delta_0 = \frac{\delta\rho}{\overline{\rho}} = \frac{\rho-\overline{\rho}}{\overline{\rho}} = \frac{M/V-\overline{\rho}}{\overline{\rho}} = (2\pi)^3\Big(\frac{\eta}{v}\Big)-1\ , \tag{10}$$

where $\overline{\rho} = \overline{n}\,E = E_0(\lambda m/g)^3/(2\pi)^3$ denotes the average background energy density of oscillons. The η prefactor can be restated in terms of $\overline{n}$ as $(\overline{\rho}V/E_0) \simeq (\lambda m/g)^3V/(2\pi)^3 = \beta^{-3/5}\overline{n}V$. The probability distribution of δ_0 is then

$$P_{\delta_0}(\delta_0) = P_\eta(\eta)\Big|\frac{d\eta}{d\delta_0}\Big| = \frac{v}{(2\pi)^3}P_\eta\Big(\frac{v}{(2\pi)^3}[1+\delta_0]\Big)\ . \tag{11}$$

The initial overdensities evolve due to gravitational self-attraction and grow according to the scale factor during the oscillon matter-dominated era as $\delta(t) = \delta_0 a(t) = \delta_0(t/t_0)^{2/3}$. Once overdensities δ exceed the critical threshold $\delta_c \sim 1$, regions start collapsing and forming black holes. Using Eq. (10) to exchange δ_c for η_c, we obtain the condition for η to be super-critical

$$\eta \geq \eta_c(a,V) = \frac{\overline{\rho}V}{E_0}\left(\frac{\delta_c}{a}+1\right)\ . \tag{12}$$

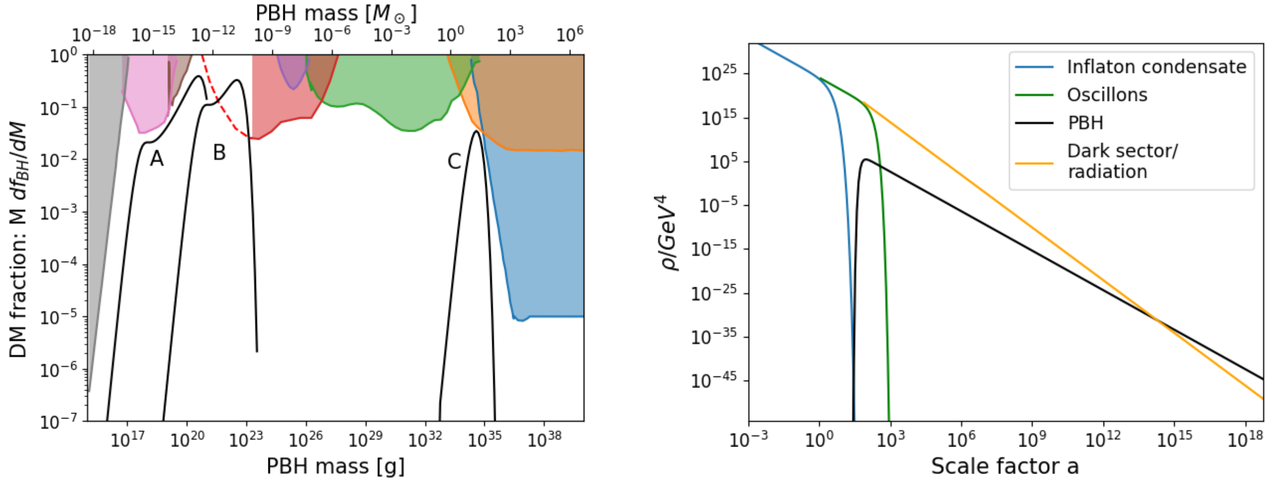

FIG. 2: [Left] DM fraction in primordial black holes. Fits for parameter choices corresponding to Model A, B, C are shown. Constraints from extragalactic γ-rays from BH evaporation [52] (EGγ), femto-lensing [53] (FL), white dwarf abundance [54] (WD), Kepler star milli/micro- lensing [55] (K), Subaru HSC micro-lensing [56] (HSC) and MACHO/EROS/ OGLE micro-lensing [57] (ML) are displayed. Dashed line indicates that HSC constraints [58] are expected to be weaker than reported when PBH Schwarzschild radius becomes smaller than the wavelength of light [59, 60]. [Right] Evolution of cosmological density for parameters of Model A. Contributions from the un-fragmented inflaton, oscillons, PBHs as well as the radiation sector are shown.

The total fraction of super-critical overdensities as a function of scale factor a can be found through integration

$$P(\delta \geq \delta_c) = \int_{\delta_c=1}^{\infty} \frac{d\delta}{a} P_{\delta_0}(\delta/a). \tag{13}$$

In Figure 1 we display this fraction for several different volume values. As can be seen, already within a few a from the time of fragmentation the amount of super-critical regions becomes significant.

Not all super-critical regions result in a black hole. Unlike the radiation-era PBH formation [61], absence of pressure gradient in the matter-dominated era greatly enhances black hole production [33]. On the other hand, non-spherical density anisotropies now play a dominant role with final stages of collapse being described by a "Zel'dovich pancake" [62], whose parameter distribution can found in [63]. Using Thorne's hoop conjecture [64] as a requirement for formation of the black hole horizon, PBH production was recently re-analyzed in [65] (see discussion in text for comparison with [33]). The probability for a super-critical overdensity region to result in a black hole is given by

$$B(M) \simeq 0.05556\, \delta^5 \left(\frac{M}{\overline{M}(V_H)}\right)^{10/3}, \tag{14}$$

where $\overline{M}(V_H) = (4\pi/3)\overline{\rho}/H_i^3$ denotes the average mass in the Hubble horizon volume $V_H(t) = (4\pi/3)t^3$ at fragmentation, with $t = 1/H$. We have further checked that including the effects of PBH spins [66], relevant for small overdensities, will not significantly alter our results.

At some scale factor a_R the oscillon matter-dominated era ends and the Universe is reheated, entering the radiation-dominated phase. Specifying the energy density of the overdensities as $\rho = M/V = E_0\eta/V$, the PBH spectrum at a_R is given by

$$\frac{d\langle\rho_{\rm PBH}\rangle}{d\eta} = \int_{V_{\rm min}}^{V_{\rm max}} \frac{dV}{V} \left(B(\eta)\frac{E_0\eta}{V}\right) P_\eta(\eta) \times \theta[\eta - \eta_c(a_R, V)]\theta\left[\rho_0 - \frac{E_0\eta}{V}\right], \tag{15}$$

where $\theta[x]$ is the Heaviside step function and $V_{\rm min}$, $V_{\rm max}$ are the average volume of a single oscillon and the Hubble horizon volume, respectively. The first step function selects super-critical regions. The second step function imposes energy conservation by requiring that an overdensity doesn't exceed the inflaton energy density, assuming that both have the same volume. The inflaton energy density ρ_0 at the bottom of its potential can be found from the mass term $\rho_0 a_i^3 = (1/2)m^2\overline{\phi}_i$, where $\overline{\phi}_i = \sqrt{3\lambda/5g^2}m$. A similar relation can be obtained directly from the Friedmann equations. In order to get the present day distribution we must redshift the results obtained at fragmentation time. The redshift factor $(a_F/a_R) = g_{*S}^{1/3}(T_F)T_F/g_{*S}^{1/3}(T_R)T_R$ accounts for evolution from T_R (defined by $\rho_R(T_R) = (\pi^2/30)g_*(T_R)T_R^4$) to $T_0 = 2.7\text{K} = 2.3$ meV. Here, g_* denotes the relevant number of relativistic degrees of freedom. Reverting from η

Model	β	$(\lambda/g)^2$	a_R	m_ϕ (GeV)	$\overline{\phi}_{\rm frag}$ (GeV)	H_i (GeV)	Γ_ϕ (GeV)	$T_{\rm R}$ (GeV)	$f_{\rm PBH}$
A	56.1	0.2	70	5×10^{-5}	6.7×10^{15}	4.6×10^{-7}	1.2×10^{-10}	6.6×10^{3}	1.0
B	35.5	0.2	20	9×10^{-8}	1.1×10^{16}	1.3×10^{-9}	2.2×10^{-12}	9.0×10^{2}	1.0
C	10.7	0.2	2	1×10^{-20}	5.3×10^{16}	7.2×10^{-22}	3.8×10^{-23}	3.8×10^{-3}	5.1×10^{-2}

TABLE I: Parameter sets for three specific model realizations (Model A, B, C). In Models A, B PBHs can account for all of the DM, while model C allows for PBHs to contribute to the observed LIGO black hole merger events. Vertical double line divides the input quantities [left-side] β, $(\lambda/g)^2$, reheating time $a_R = a(T_R)$, inflaton mass m_ϕ and the derived quantities [right-side]: inflaton VEV at fragmentation $\overline{\phi}_{\rm frag}$, initial Hubble rate H_i, inflaton decay rate Γ_ϕ, reheating temperature T_R and the fraction of DM in PBHs $f_{\rm PBH}$.

back to M, the fraction of DM residing in PBHs is then

$$\frac{df_{\rm PBH}}{dM} = \frac{1}{\rho_{\rm DM} a^3}\frac{d\langle\rho_{\rm PBH}\rangle}{dM}\ , \tag{16}$$

where $\rho_{\rm DM}$ is the present-day DM density and the $1/a^3$ factor accounts for the redshift In Figure 2 we display the fraction of PBHs as DM for several specific parameter sets (denoted as “Model A”, “B”, “C”) along with the current experimental constraints. Model A, B correspond to the region where PBHs can make up all of the DM, while Model C covers the region where PBHs can contribute to the observed LIGO black hole merger events [23, 26, 27]. Exact values of the parameters, including both the input and the derived quantities, can be found in Table I. We note that Model C is phenomenologically not viable; the relevant parameters are shown for completeness.

In Figure 2 we display the cosmological history of the setup, showing energy density evolution of the inflaton, oscillons, PBHs as well as radiation from reheating. During inflation, the inflaton dominates the Universe. As the inflaton settles at the bottom of potential the Universe becomes matter-dominated with the density scaling as a^{-3}. After fragmentation of the inflaton into oscillons and PBH formation the Universe is reheated, becoming radiation-dominated with the density scaling as a^{-4}. At redshift $z \approx 3600$ dark matter in the form of PBHs, whose density scales also as a^{-3}, overtakes the radiation contribution, and the Universe again enters matter-dominated regime. Unlike the case of Q-balls [8, 34], there is no intermediate radiation-dominated era before the fragmentation time in our setup, since oscillons form directly from the inflaton during the early stages of reheating.

We further comment on two important cosmological aspects of the setting, the inflationary phase and reheating. Taken at face value, the potential of Eq. (1) produces unphysical perturbation spectrum during inflation due to the dominance of ϕ^6 term (see [41] for discussion). However, the field value $\overline{\phi}$ at the bottom of potential is far below the Planck scale that sets the initial inflaton displacement. Hence, our region of interest where inflaton oscillates near the potential minimum is decoupled from the large values that determine the inflationary phase. In this work we remain agnostic regarding the exact shape of the potential and the early Universe dynamics. We focus on the effective potential at a relatively small vacuum expectation value (VEV), well below the values that are relevant for structure formation.

The Universe is reheated from oscillon decay. If PBHs are to constitute a significant fraction of DM, the inflaton must be very light, and the reheating temperature is very low (see Table I). To avoid affecting the Big Bang nucleosynthesis (BBN), reheating should occur above $T \gtrsim 4$ MeV scale (e.g. [67]). While it is commonly assumed that the Universe was reheated to a much higher temperature, a cosmological history with $\sim$MeV reheating is possible and is consistent with observations [68, 69]. Neglecting the oscillon quantum decay [70], the allowed direct inflaton decay channels are limited by the total invariant mass. The simplest decay mode is to photons $\phi \to \gamma\gamma$, proceeding through an effective $(g_{\gamma\gamma}/4)F_{\mu\nu}F^{\mu\nu}\phi$ operator, where $F_{\mu\nu}$ is the electromagnetic field strength tensor and $g_{\gamma\gamma}$ is the coupling. The relevant decay rate is given by $\Gamma_{\phi\to\gamma\gamma} = (g_{\gamma\gamma}^2/64\pi)m_\phi^3$. However, axion-like particle searches already strongly constrain this channel [71, 72]. Thus, without an extended dark sector, Model C is not viable. In a more complicated model, the reheating may be possible if the inflaton decays into some dark sector particles, which produce the Standard Model degrees of freedom via mixing (e.g. [73, 74]). Generating the matter-antimatter asymmetry in a low-reheating scenario also presents a model building challenge.

In summary, inflaton fragmentation into oscillons can lead to formation of primordial black holes in a single-field inflation model or other models that admit oscillon solutions. This novel production mechanism can generate a sufficient density of PBHs to account for all or part of dark matter. It is also possible that solar-mass black holes can be produced this way, but the required mass of the inflaton is very small, and the need for reheating and baryogenesis will lead to more complicated models.

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* ecotner@physics.ucla.edu
† kusenko@ucla.edu
‡ vtakhist@physics.ucla.edu

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