

Regurgitated dark matter

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(Received 17 September 2023; accepted 6 August 2024; published 20 September 2024)

We present a new paradigm for the production of the dark matter (DM) relic abundance based on the evaporation of early Universe primordial black holes (PBHs) themselves formed from DM particles. As a concrete realization, we consider a minimal model of the dark sector in which a first-order phase transition results in the formation of Fermiball remnants that collapse to PBHs, which then emit DM particles. We show that the regurgitated DM scenario allows for DM in the mass range $\sim 1\text{--}10^{16}$ GeV, thereby unlocking parameter space considered excluded.

DOI: [10.1103/PhysRevD.110.L051702](https://doi.org/10.1103/PhysRevD.110.L051702)

Introduction. Dark matter (DM) constitutes $\sim 85\%$ of all matter in the Universe, as determined by a multitude of astronomical observations (for reviews, see, e.g., [1,2]). Despite decades of efforts to detect its nongravitational interactions, the nature of DM remains mysterious.

A significant focus has been on weakly interacting massive particles (WIMPs) as the DM paradigm [3,4], with typical masses in the GeV to multi-TeV range that often appear in theories that address the hierarchy problem of the standard model (SM). The scenario of decoupling from the thermal bath in the early Universe, which successfully explains cosmological observations such as the light element abundances [5], suggests that WIMPs with typical electroweak-scale masses and annihilation cross sections can readily account for the observed DM relic abundance through thermal freeze-out (the so-called “WIMP miracle”). Sensitive experimental searches (e.g., [6,7]) significantly

constrain the parameter space of minimal WIMP scenarios. DM scenarios based on additional or number-changing DM particle interactions, such as the strongly interacting massive particle miracle [8], provide alternative approaches to achieving the DM relic abundance.

In this work, we present a novel paradigm, which we call *regurgitated dark matter* (RDM), for producing the DM relic abundance. It is based on the evaporation of early Universe primordial black holes (PBHs), themselves constituted by DM particles. PBHs scramble and reemit DM particles with altered energy momentum and abundance distributions, distinct from that of the original DM particles in the thermal bath that formed the PBHs; i.e., these properties are not determined by direct DM interactions as in conventional DM production mechanisms. While particle DM emission from evaporating PBHs has been studied (e.g., [9–11]), in the RDM scenario, PBHs originate from the same DM particles that later constitute the DM relic abundance and not from some distinct mechanism; for PBH production mechanisms, see, e.g., [12–25]. As we demonstrate, RDM can open a new window in the WIMP parameter space with either fermion or scalar DM particles.

Model. We illustrate an elegant realization of RDM in the context of a first-order phase transition (FOPT) in an asymmetric dark sector that produces Fermiball remnants composed of dark sector particles that subsequently collapse

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to PBHs. The dark sector particles that form the PBHs are emitted by these PBHs through Hawking evaporation. Depending on the particle mass and PBH mass, RDM particles may be relativistic at the epoch of big bang nucleosynthesis (BBN) or contribute to warm DM. We note, however, that our RDM mechanism is general and may be realized in the context of other scenarios, such as the collapse of solitonic macroscopic objects to PBHs or in models with additional dark sector forces and SM portals beyond the Higgs portal. We leave the exploration of such possibilities for future work.

We consider the model of Refs. [26–30] given by

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - \frac{\mu^2}{2} \phi^2 - \frac{\kappa}{2} \phi^2 (\mathcal{H}^\dagger \mathcal{H}) - V(\phi) + \bar{\chi} i \not{\partial} \chi - y_\chi \phi \bar{\chi} \chi, \quad (1)$$

where $\mathcal{L}_{\text{SM}}$ is the standard model (SM) Lagrangian, dark sector fermions χ , $\bar{\chi}$ and scalar ϕ interact via an attractive Yukawa force with coupling y_χ , and ϕ interacts with the SM sector through the Higgs $\mathcal{H}$ doublet portal coupling κ . On receiving thermal corrections, the potential $V(\phi)$ becomes $V(\phi, T)$, which triggers a FOPT below the critical temperature, forming Fermiball remnants that collapse to PBHs. We remain agnostic about the details of the potential, allowing a general discussion of RDM. We assume the existence of an asymmetry between the number density of χ and $\bar{\chi}$ with $\eta_\chi = (n_\chi - n_{\bar{\chi}})/s(T_\star)$, where $s(T_\star) = 2\pi^2 g(T_\star) T_\star^3/45$ is the entropy number density with relativistic degrees of freedom (d.o.f.) $g(T)$ at temperature $T_\star$ of the FOPT; we neglect the small contribution of the dark sector d.o.f. to $g(T)$. Such an asymmetry can be realized in a variety of asymmetric DM mechanisms [31–33], with Fermi-degenerate remnants dominated by χ .

Formation of fermiballs. We consider the following thermal history of cosmology for production of RDM, illustrated schematically in Fig. 1.

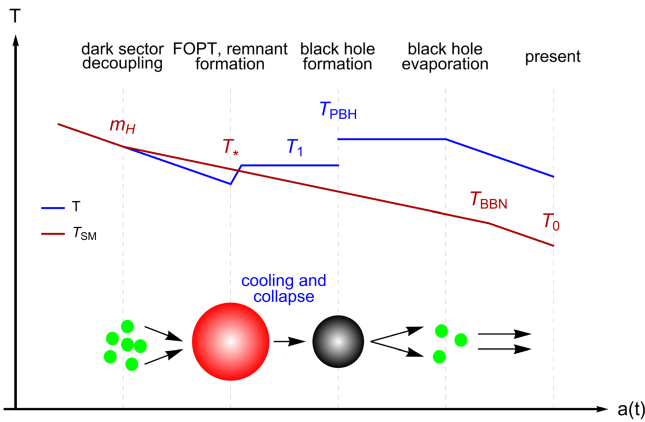


FIG. 1. Cosmological thermal history of RDM production. The dark sector particles in the Fermiball are reemitted at a higher temperature after black hole formation.

At first, the dark sector and SM particles are in thermal equilibrium after inflationary reheating, which can occur either due to the Higgs portal coupling κ or inflaton sector couplings. As the Universe expands and temperature decreases below the electroweak phase transition at $T \sim 160$ GeV, the dark sector decouples from the visible sector at $T \sim m_h = 125$ GeV due to the diminishing interactions between ϕ and the Higgs. Hereafter, the SM and dark sector temperatures T_{SM} and T , respectively, evolve separately. The effective number of relativistic degrees of freedom (d.o.f.) of the SM at dark sector decoupling is $g(T_{\text{SM}}^{\text{dec}})$.

At a temperature $T_\star$, a FOPT is triggered by the dark scalar potential $V(\phi, T_\star)$. The phase transition changes the expectation value of ϕ , from $\langle \phi \rangle = 0$ in the false vacuum to $\langle \phi \rangle = v_\star$ in the true vacuum. The FOPT proceeds through bubble nucleation (with expanding bubble wall speed v_w) and can be characterized by the following parameters [34]: β is the inverse duration of the FOPT, and $\alpha_D \simeq 30\Delta V/\pi^2 g_D T_\star^4$ quantifies its strength. Here, $g_D = 4.5$ is the d.o.f. of the dark sector, and ΔV is the potential energy difference between the false and true vacua.

The FOPT can readily induce a significant mass gap $\Delta m_\phi, \Delta m_\chi \gg T_\star$ between the vacua, with massless fermions acquiring mass $m_\chi = y_\chi v_\star$ via the Yukawa coupling. For the particles to be trapped, the dark sector particle masses in the true vacuum must exceed the FOPT temperature:

$$m_\chi \gg T_\star, \quad m_\phi = \left(\frac{\partial^2 V(\phi, T_\star)}{\partial \phi^2} \right)^{1/2} \Big|_{\phi=v_\star} \gg T_\star. \quad (2)$$

This condition can be fulfilled in supercooled scenarios with large $v_\star/T_\star$ [35–39], or with $y_\chi \gg 1$ [40,41]. Then the mass of dark sector particles is significantly larger than their thermal kinetic energy and cannot penetrate into the true vacuum bubbles. As true vacuum bubbles expand, dark sector particles are efficiently trapped in contracting regions of false vacuum. In much of the parameter space, explicit calculations confirm that the trapping fraction of dark sector particles in the false vacuum remnants is ~ 1 if Eq. (2) is satisfied. The remnants get compressed to form nontopological solitonic Fermiball remnants [26–30,42,43]. We assume that the number of d.o.f. does not significantly vary between decoupling and the FOPT, which will not affect our conclusions.

Fermiball cooling. The dark sector temperature of the Fermiballs will be $T_1 = (90\Delta V/\pi^2 g_D)^{1/4}$ during the slow remnant cooling process [29]. The Fermiballs cool via SM particle production through the Higgs portal. A detailed analysis of Fermiball cooling for our regimes of interest can

be found in the Supplemental Material [44]. The asymmetry η_χ ensures that a population of χ particles survives after annihilation to ϕ 's. As the Fermiball cools, these particles dominate until the Fermiball collapses into a black hole.

The dominant cooling channel for Fermiballs depends on the dark sector temperature T_1 . For T_1 below the electroweak scale, the cooling rate is suppressed by the Higgs mass. For small values of κ , we have verified that volumetric cooling occurs with a rate $\dot{C} = n^2 \langle 2E \rangle \sigma v_{\text{rel}}$; see the Supplemental Material [44] for details. The scalar ϕ follows a thermal Bose-Einstein distribution with number density $n = (\zeta(3)/\pi^2) T^3$ and average energy $\langle E \rangle \simeq 2.7 T_1$.

In this regime, Fermiball remnants will predominantly cool through $\phi\phi \rightarrow f\bar{f}$ emission, where f is the heaviest available fermion with mass m_f . The remnant is initially supported by thermal pressure, but as it cools, it becomes supported by the Fermi degeneracy pressure of the asymmetric χ population. In the volumetric cooling regime, this transition happens at a temperature,

$$T_{\text{SM}}^{\text{tr,low}} \simeq (10^4 \text{ GeV}) \kappa \left(\frac{T_1}{1 \text{ GeV}} \right)^{3/2} \frac{m_f}{1.27 \text{ GeV}} \times \left(\frac{g_D}{4.5} \right)^{-1/2} \left(\frac{g(T_{\text{SM}}^{\text{tr}})}{106.75} \right)^{-1/4}. \quad (3)$$

For T_1 above the electroweak scale, direct Higgs production, $\phi\phi \rightarrow \mathcal{H}\mathcal{H}$, can occur. For small κ , we have volumetric cooling with the transition temperature,

$$T_{\text{SM}}^{\text{tr,high}} \simeq (10^4 \text{ TeV}) \kappa \left(\frac{T_1}{1 \text{ TeV}} \right)^{1/2} \left(\frac{g_D}{4.5} \right)^{-1/2}. \quad (4)$$

For larger κ in both temperature ranges, the mean free path of the particles can become shorter than the Fermiball radius, and blackbody surface cooling dominates. In this case, the transition time is negligible, i.e., $T_{\text{SM}}^{\text{tr}} \sim T_\star$; see the Supplemental Material [44] for details. Hence, the resulting DM abundance is determined primarily by the black hole evaporation timescale.

Black hole formation. Black holes are formed from the cooling Fermiballs when the length scale associated with attractive Yukawa force $\sim 1/m_\phi$ is comparable with the mean separation of χ ($\bar{\chi}$) inside. In practice, this collapse occurs shortly after the transition to Fermi degeneracy pressure, at temperature $T_{\text{SM}}^{\text{tr}}$ [29,30]. (Note that gravitational collapse is also possible [46].) This instability is ensured for $\alpha_D > 0.01$ [30], which is readily satisfied for $\alpha_D > 1/3$ so that the initial remnant shrinks efficiently. Then, the average mass of PBHs formed from the collapse of Fermiballs is [29]

$$\bar{M}_{\text{PBH}} \simeq (6.61 \times 10^{14} \text{ g}) \alpha_D^{1/4} \left(\frac{v_w}{0.7} \right)^3 \frac{\eta_\chi}{10^{-10}} \left(\frac{\beta/H}{1000} \right)^{-3} \times \left(\frac{g(T_\star)}{g(T_{\text{SM}}^{\text{dec}})} \right)^{-2/3} \left(\frac{T_\star}{1 \text{ GeV}} \right)^{-2}. \quad (5)$$

The number density of these PBHs is

$$n_{\text{PBH}}(T_{\text{SM}}) = 3.83 \times 10^8 \left(\frac{v_w}{0.7} \right)^{-3} \left(\frac{\beta/H}{1000} \right)^3 \times \frac{g_s(T_{\text{SM}}) g(T_\star)^{1/2} T_{\text{SM}}^3 T_\star^3}{M_{\text{pl}}^3}, \quad (6)$$

with $g_s(T_{\text{SM}})$ the number of entropic d.o.f. of the visible sector.

PBHs will eventually evaporate through Hawking radiation [47] at a time t_{PBH} after formation, as reviewed in the Supplemental Material [44]. If PBHs come to dominate the matter density, then evaporation will reheat the Universe to a temperature,

$$T_{\text{SM}}^{\text{RH}} = 50.5 \text{ MeV} \left[\frac{\bar{M}_{\text{PBH}}}{10^8 \text{ g}} \right]^{-\frac{1}{2}} \left[\frac{g(T_{\text{SM}}^{\text{RH}})}{10} \right]^{-\frac{1}{4}} \left[\frac{g_{\text{H,SM}}}{108} \right]^{\frac{1}{2}}, \quad (7)$$

where $g_{\text{H,SM}}$ is the number of Hawking d.o.f. for the SM sector, which we elaborate on below. If there is no PBH-dominated era, then the prefactor in Eq. (7) becomes $T_{\text{SM}}^{\text{evap}} \sim 43 \text{ MeV}$. The time at which evaporation occurs is the sum of the three timescales: the PBH lifetime, the formation (cooling) time, and $t_\star$. Because of the hierarchy of scales, the evaporation occurs at $\sim \min(T_{\text{SM}}^{\text{tr}}, T_{\text{SM}}^{\text{RH}}, T_\star)$.

There are several key timescales in our setup. We are primarily interested in PBH evaporating into sufficient quantities of RDM before BBN to maintain the successful predictions of the light element abundances. This sequence includes the production of Fermiballs, their subsequent cooling and collapse into PBHs, and PBH evaporation. The relevant timescales for Fermiball formation and PBH collapse are related to the temperatures $T_\star$ and $T_{\text{SM}}^{\text{tr}}$, respectively, and the evaporation temperature is $T_{\text{SM}}^{\text{RH}}$ ($T_{\text{SM}}^{\text{evap}}$). To evade BBN constraints, we require that $T_\star, T_{\text{SM}}^{\text{tr}}, T_{\text{SM}}^{\text{RH}} \geq 5 \text{ MeV}$. The condition for PBH domination before they evaporate is $\bar{M}_{\text{PBH}} n_{\text{PBH}}(T_{\text{SM}}^{\text{RH}}) > \rho_{\text{SM}}$.

Regurgitated dark matter. The dark sector particles are regurgitated during PBH evaporation and can constitute the main component of DM either in the WIMP mass range, $1 \text{ GeV} \lesssim m_\chi \lesssim 1 \text{ TeV}$, or the superheavy mass range with $m_\chi \gg 1 \text{ TeV}$. The energy density of RDM depends on the PBH mass fraction at evaporation. If the PBHs evaporate during the radiation dominated era, the density of regurgitated particles is suppressed by $\rho_{\text{PBH}}/(\rho_{\text{PBH}} + \rho_{\text{SM}}) \simeq \rho_{\text{PBH}}/\rho_{\text{SM}}$. The ratio is unity if PBHs dominate the matter density.

The relative Hawking emission rates of dark sector particles to SM particles is given by the ratio of their effective Hawking d.o.f. $g_{H,D}/g_{H,SM}$, with $g_{H,D} = 5.82$ for the dark sector and $g_{H,SM} = 108$ for the SM sector [48,49]; see the Supplemental Material [44]. Most emitted particles have masses smaller than the Hawking temperature T_{PBH} of the PBH. So, $T_{PBH} \gtrsim m_\chi, m_\phi \gg T_1, T_\star$ holds for dark sector particle emission. While both ϕ and χ are emitted, the scalar ϕ develops a mixing with the SM Higgs after the electroweak and dark sector phase transitions, allowing for the decay of ϕ into SM particles. In the relevant regions of allowed parameter space, this decay timescale is shorter than the lifetime of the Universe. Hence, in this particular model realization of the RDM paradigm, ϕ does not constitute a significant DM component. In the following, we focus on the abundance of the stable fermion χ . Although ϕ is unstable and does not contribute to the current mass density, its parameters can still be constrained by considerations of the cooling rate and the Higgs invisible decay. In these instances, we assume $m_\phi = m_\chi$.

When the initial Hawking temperature is smaller than the particle mass, the dark sector particles are emitted once the Hawking temperature reaches their mass in the true vacuum, $\epsilon_{em} T_{PBH} \simeq m_\chi$ where $\epsilon_{em} = 2.66, 4.53, 6.04$ for spin $s = 0, 1/2, 1$ particles, respectively [50]. This higher temperature corresponds to a lighter PBH mass below which these heavy particles can be emitted:

$$M_{PBH}^{em} = (1.06 \times 10^8 \text{ g}) \epsilon_{em} \left(\frac{m_\chi}{10^5 \text{ GeV}} \right)^{-1}. \quad (8)$$

The bulk of emitted particles will be nonrelativistic, with the initial density at emission reduced by a factor of $(M_{PBH}^{em}/\bar{M}_{PBH})$, such that $\rho_\chi/\rho_{SM} = (M_{PBH}^{em}/\bar{M}_{PBH})(g_{H,\chi}/g_{H,SM})$.

From the emission spectrum of PBHs, the average energy for heavy particles is $\bar{E} = 2m_\chi$. Thus, the present day DM mass density of such primarily nonrelativistic emitted DM particles is

$$\begin{aligned} \Omega_{DM,\chi}^{NR} &= \frac{\Omega_{r,0} g_s^{1/3} (T_{SM}^{RH}) T_{SM}^{RH} M_{PBH}^{em} g_{H,\chi}}{2 g_s^{1/3} (T_{SM}^0) T_{SM}^0 \bar{M}_{PBH} g_{H,SM}} \\ &= 1.24 \times 10^5 \epsilon_{em} g_{H,\chi} \left(\frac{m_{(\phi,\chi)}}{10^5 \text{ GeV}} \right)^{-1} \\ &\quad \times \left(\frac{\bar{M}_{PBH}}{10^8 \text{ g}} \right)^{-5/2} \left(\frac{g(T_{SM}^{RH})}{10} \right)^{1/12}, \end{aligned} \quad (9)$$

where $T_{SM}^0 = 2.73 \text{ K}$ and $g_s(T_{SM}^0) = 3.9$ are the present day values. An additional factor of ρ_{PBH}/ρ_{SM} appears if the PBHs are subdominant (see Supplemental Material [44]).

If the particles are emitted relativistically, the present day mass density contribution is reduced by the redshifting of these particles until they become nonrelativistic. The bulk

of the Hawking radiation emitted particles will have a Lorentz boost factor $\gamma \simeq \epsilon_{em} T_{PBH}/m_\chi$. The resulting density of these initially relativistic dark sector particles is

$$\Omega_{DM,\chi}^R = \Omega_{DM,\chi}^{NR} \times \frac{4}{3} \frac{m_\chi}{\epsilon_{em} T_{PBH}} \frac{\bar{M}_{PBH}}{M_{PBH}^{em}}. \quad (10)$$

The behavior is opposite to the nonrelativistic case, with lighter particles being less dense.

Dark matter detection. As dark matter, χ can be observed in direct detection experiments by interactions through the Higgs portal coupling in Eq. (1). The resulting elastic scattering cross section of χ on nucleons N is given by [51,52] (see Supplemental Material [44] for details)

$$\begin{aligned} \sigma_{\chi N} &= \frac{\kappa^2}{\pi m_h^4} \left(\frac{m_\chi^2}{m_\phi^2 - m_h^2} \right)^2 \frac{m_N^4 f_N^2}{(m_\chi + m_N)^2} \\ &\sim \frac{3.5 \times 10^{-38} \text{ cm}^2 \times \kappa^2}{\left(\frac{m_\chi}{1 \text{ GeV}} + 1 \right)^2} \times \left(\frac{m_\chi^2}{m_\phi^2 - m_h^2} \right)^2, \end{aligned} \quad (11)$$

where $m_N \simeq 1 \text{ GeV}$ is the nucleon mass, and $f_N \sim 0.3$ is Higgs-nucleon interaction parameter.

Employing Eq. (11), we recast existing bounds on the spin-independent scattering cross section into constraints on the coupling κ . In Fig. 2, we display constraints on κ (upper shaded region) as well as predictions (solid colored lines) for χ to constitute all of the DM abundance for different values of β/H and η_χ . In the lower shaded region labeled ‘‘Cool,’’ PBHs form after BBN via Fermiball cooling. Clearly, the regurgitated χ can saturate the DM relic abundance for a wide range of masses. The final abundance has a κ dependence if the cooling time is longer than the lifetime of the PBH and the transition time $t_\star$. In the WIMP mass range, the cooling rate depends on fermion channels that open sequentially with increasing $m_\chi = 20T_1$. Here, $g(T_\star)$ also changes similarly and affects the PBH number density through Eq. (6).

The Yukawa interaction between χ and ϕ keeps the χ in thermal equilibrium after it decouples from the SM. Therefore, we calculate χ 's freeze-out abundance by assuming the two particles freeze-out together when ϕ decouples from the SM plasma. We interpret the results for thermal freeze-out production of Ref. [54] for $\phi\phi \rightarrow f\bar{f}$ in the nonrelativistic limit (see Supplemental Material [44]) and plot the result as the solid black curve in Fig. 2. While WIMP masses are strongly constrained, RDM can be efficiently produced in the unconstrained parameter space below $m_\chi \sim \text{TeV}$.

For $m_\chi < m_h/2$, stringent bounds on κ arise from the invisible Higgs decay branching fraction to DM particles $\text{Br}(H \rightarrow \text{inv})$ constrained by Large Hadron Collider (LHC) data at 95% confidence level [51]. We show this bound in

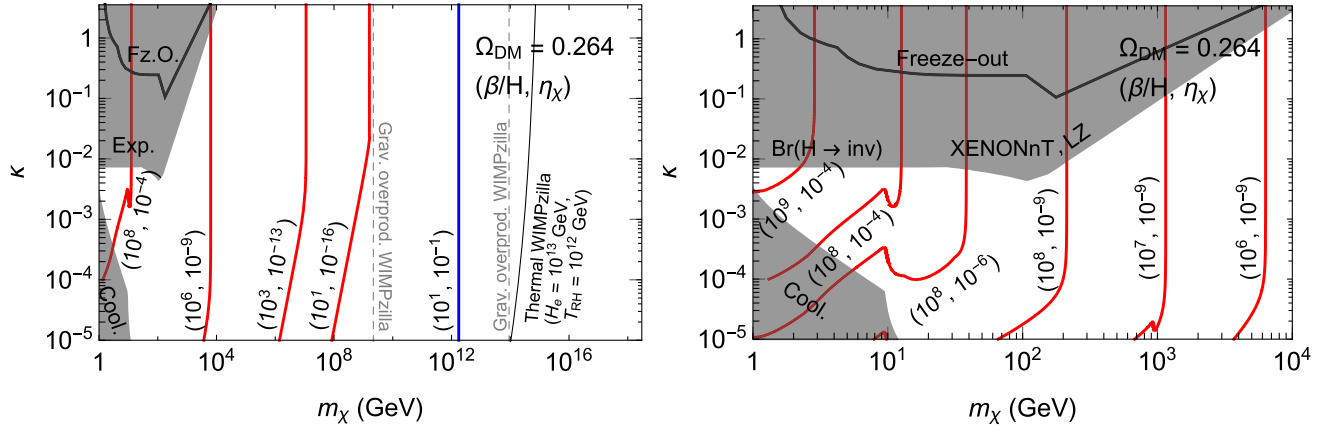


FIG. 2. Constraints on the Higgs portal coupling κ as a function of m_χ (assuming $m_\phi = m_\chi$) over a wide mass range [left] and for the WIMP mass range [right]. On the solid iso- $\Omega_{\text{DM}} = 0.264$ contours, χ constitutes the entire DM abundance for specific choices of the inverse duration of the FOPT β/H and dark sector fermion asymmetry η_χ . The PBHs dominate the Universe's energy density on the blue contours and are subdominant on the red contours. We fix $\alpha_D = 1$, $v_w = 0.67$, $g_{D,*} = 4$ and $T_1 \simeq T_* \simeq 0.05 m_\chi$. Also displayed are the most restrictive bounds from several experiments (labeled “Exp.” in the left panel), including the 95% CL bound from the invisible Higgs decay branching fraction [51], and the 90% CL bounds from XENONnT [6] and LUX-ZEPLIN (LZ) [53]. The solid black curve corresponds to conventional Higgs portal WIMP freeze-out, and the solid and dashed gray curves correspond to thermal production and gravitational overproduction of superheavy WIMPZillas, respectively. In the lower shaded region labeled “Cool.,” PBHs form after BBN.

Fig. 2 with $m_\phi = m_\chi$. Direct detection constraints on κ weaken for heavier m_χ since $\sigma_{\chi N} \propto 1/m_\chi^2$ for $m_\chi \sim m_\phi \gg m_h \gg m_N$.

In the mass range $10^9 \text{ GeV} \lesssim m_\chi \lesssim 10^{14} \text{ GeV}$, gravitational overproduction of WIMPZillas [55–58] can restrict χ from being a viable DM candidate depending on the Hubble rate H_e at the end of inflation [59]; see the gray dashed lines in Fig. 2. However, if inflation occurs at a lower energy scale, the abundance of WIMPZillas can be significantly suppressed.

If the WIMPZilla has additional interactions, a thermal relic can be realized. For the Higgs portal scenario [59], we show the case of $H_e = 10^{13} \text{ GeV}$ with the reheating temperature $T_{\text{SM}}^{\text{RH}} = 10^{12} \text{ GeV}$. For even heavier DM masses, $m_\chi \gtrsim 3 \times 10^{16} \text{ GeV}$, stringent bounds originate from DEAP-3600 and mica searches [60], but they fall in a region with no reliable scaling relation. Furthermore, a robust theoretical bound for pointlike DM [61] cannot be meaningfully applied to κ since this would imply that the nuclear scattering cross section of χ plateaus to a maximum value even for an arbitrarily large κ .

Gravitational waves (GWs) from the FOPT can provide a correlated signature with DM detection in scattering experiments. While detailed predictions depend on specifics of the FOPT, the peak GW frequency is expected to be $f_{\text{GW}} \simeq \mathcal{O}(10^{-2}) \text{ mHz}(\beta/H_*/100)(T_*/1 \text{ GeV})$, which could fall in the range of upcoming interferometers such as LISA [62], Einstein Telescope [63], Cosmic Explorer [64], Big Bang Observer [65], and DECIGO [66]. Moreover, if PBHs dominate the matter density, their evaporation can

lead to induced GWs [67–71]. We leave the exploration of the associated GW signal for future work.

Conclusions. We proposed a novel paradigm of regurgitated DM, stemming from the emission of evaporating PBHs formed from the DM particles themselves. This is distinct from conventional particle DM production mechanisms since the resulting DM relic abundance is not determined by particle interactions. Intriguingly, as we demonstrate with a concrete realization, this paradigm can produce the inferred abundance of DM in a very broad mass range $\sim 1\text{--}10^{16} \text{ GeV}$ and opens up parameter space previously thought to be excluded. A stochastic background of GWs is a possible correlated signature of the scenario.

Acknowledgments. We thank J. Arakawa, J. Jeong, S. Jung, K. Kawana, J. Kim, and K. Xie for useful discussions. T.H.K. is supported by a KIAS Individual Grant No. PG095201 at Korea Institute for Advanced Study and National Research Foundation of Korea under Grant No. NRF-2019R1C1C1010050. P.L. is supported by Grant Korea NRF2019R1C1C1010050. D.M. is supported in part by the U.S. Department of Energy under Grant No. DE-SC0010504. V.T. acknowledges support by the World Premier International Research Center Initiative (WPI), MEXT, Japan and JSPS KAKENHI Grant No. 23K13109. This work was performed in part at the Aspen Center for Physics, which is supported by the National Science Foundation Grant PHY-2210452.

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