Little Red Dots and Supermassive Black Hole Seed Formation in Ultralight Dark Matter Halos

Dongsu Bak*
Physics Department, University of Seoul, Seoul, 02504, Korea

Jae-Weon Lee†
Department of Electrical and Electronic Engineering, Jungwon University, 85 Munmu-ro, Goesan-eup, Goesan-gun, Chungcheongbuk-do, 28024, Korea.

We investigate how supermassive black hole (SMBH) seeds form in the early Universe at the centers of ultralight dark matter (ULDM) halos. Focusing on the ULDM Jeans scale, we identify the critical conditions under which high-redshift baryonic gas, strongly confined by central solitonic cores of the halos, undergoes direct and monolithic collapse. The solitonic potential naturally drives rapid inflow and shock heating, allowing the gas to exceed the critical atomic-cooling and fragmentation-suppression threshold of $\sim 3\times10^4$K without invoking an external UV background. We derive semi-analytic relations for the halo mass, soliton mass, baryonic core radius, and thermodynamic state of the gas, including the effects of baryonic contraction. These relations simultaneously determine the minimum and maximum SMBH seed masses as functions of redshift. In this framework, pristine gas clouds that satisfy the temperature threshold collapse without fragmentation, forming SMBH seeds with characteristic masses of order $\sim 10^5 M_\odot$, while systems below the threshold are expected to form compact star clusters instead. Our model also implies an upper limit on the attainable SMBH mass, predicting a maximum mass scale of order $\sim 10^{10} M_\odot$, consistent with the most massive quasars observed to date. The ULDM particle mass required to reproduce the inferred seed mass scale, $m \simeq 10^{-22}$eV, coincides with the value favored by galactic-scale observations, providing a unified explanation for the characteristic masses of both galactic cores and early SMBH seeds. Our model predicts efficient SMBH seed formation at redshifts $z \gtrsim 10$ and offers a natural interpretation of recently observed little red dots as SMBHs embedded in compact, hot, ionized gas clouds.

arXiv:2601.21676v1 [astro-ph.GA] 29 Jan 2026

I. INTRODUCTION

The detection of bright quasars less than a billion years after the Big Bang indicates that supermassive black holes (SMBHs) with masses up to $\sim 10^9 M_\odot$ already existed at the redshifts $z \gtrsim 6$, posing a serious challenge to models of black hole seeding and early growth [1, 2]. In stellar-remnant seed scenarios, maintaining nearly continuous and efficient accretion is difficult in the shallow potential wells at high redshift, which motivates massive-seed scenarios where black holes are born with initial masses of $\sim 10^4$–$10^6 M_\odot$ [1–3].

The direct-collapse black hole (DCBH) scenario assumes that a primordial, metal-free gas cloud undergoes an almost monolithic collapse to form a supermassive star or a quasi-star, which subsequently gives rise to a massive black hole seed [4, 5]. The key requirement for this scenario is the suppression of fragmentation, which could be achieved by maintaining the gas temperature at $\sim 10^4$ K and preventing

*Electronic address: dsbak@uos.ac.kr
†Electronic address: scikid@jwu.ac.kr

efficient molecular hydrogen cooling, often via a strong Lyman-Werner UV background or dynamical heating [2, 6–8]. Recent observations of little red dots (LRDs) [9] indicate that black hole seeds with masses $M_{bh} \gtrsim 10^5 M_\odot$ formed at the centers of compact ionized hot pristine gas clouds in the early Universe. It is noteworthy that the observed minimum mass of these seeds is comparable to the characteristic minimum mass of galactic cores $\sim 10^6 M_\odot$ [10]. It is therefore tempting to speculate that these two mass scales are physically connected.

One possible explanation for the characteristic mass scale of galactic cores [11] is ultralight dark matter (ULDM) [12–22], which has been extensively studied as an alternative to cold dark matter (CDM) and may alleviate small-scale issues of CDM such as the cusp problem. In this model, dark matter consists of extremely light scalar particles with mass $m \simeq 10^{-22} eV$ whose macroscopic de Broglie wavelength leads to quantum pressure supported cores (solitons) embedded in larger halos [21, 23, 24]. These solitonic cores typically have masses $M_{sol} \gtrsim 10^6 M_\odot$ and provide deep, centrally concentrated gravitational potentials that can significantly boost baryonic collapse at high redshift [25], while also helping to alleviate the final parsec problem [26].

In this paper, we explore a ULDM-assisted SMBH seed formation in which the solitonic core acts as a deep potential well for direct collapse of baryon clouds. The key idea is that baryons gravitationally confined within the soliton can undergo rapid infall accompanied by strong shock heating, thereby driving the gas into a regime where molecular cooling becomes dynamically subdominant, even without invoking an external UV background. We develop semi-analytic scalings for (i) the characteristic halo and soliton properties as a function of ULDM mass m and redshift z, and (ii) the thermodynamic state of the collapsing baryons, including the conditions under which compressional heating and shocks overwhelm H_2 cooling. We then relate these conditions to the fragmentation criterion to determine the regime in which a global collapse is expected [2–4, 27]. In this way, the ULDM soliton provides a natural, purely gravitational route to sustained high inflow rates and elevated temperatures, enabling the formation of supermassive stars with characteristic masses $\sim 10^5 M_\odot$ that can rapidly evolve into SMBH seeds consistent with high-redshift quasar and LRD observations [1, 2, 5].

Section II summarizes the ULDM quantum Jeans scales for halos and the soliton-halo relations used in this work. Section III derives the properties of baryonic core structure inside ULDM solitons, formulates the shock-heating and cooling-imbalance criteria, and presents the fragmentation suppression conditions together with the implied characteristic seed mass scale. Section IV discusses the broader implications of our model for early structure formation, including its physical interpretation of LRDs, predictions for SMBH seed formation across redshift, and the observational consequences and limitations of the scenario.

II. JEANS LENGTH AND HALO MASS FUNCTION OF ULTRALIGHT DARK MATTER

In this section, we present some scales that determine early halos and the corresponding central solitons. The quantum Jeans length scale is given by

$$\lambda_J(z) = \frac{\pi^{3/4}\hbar^{1/2}}{\left(G\rho(z)\right)^{1/4} m^{1/2}} = 9.217 \times 10^4 \,(1+z)^{-3/4} m_{22}^{-1/2} \left(\frac{0.27}{\Omega_{\rm dm}}\right)^{1/4} \left(\frac{0.7}{h}\right)^{1/2} \text{pc}, \tag{II.1}$$

where $m_{22} \equiv m/10^{-22} eV$, h is the reduced Hubble constant, and $\Omega_{\rm dm}$ is the dark matter density parameter. The dark matter density at the redshift z is

$$\rho_{dm}(z) = \rho_{dm}(0)\,(1+z)^3. \tag{II.2}$$

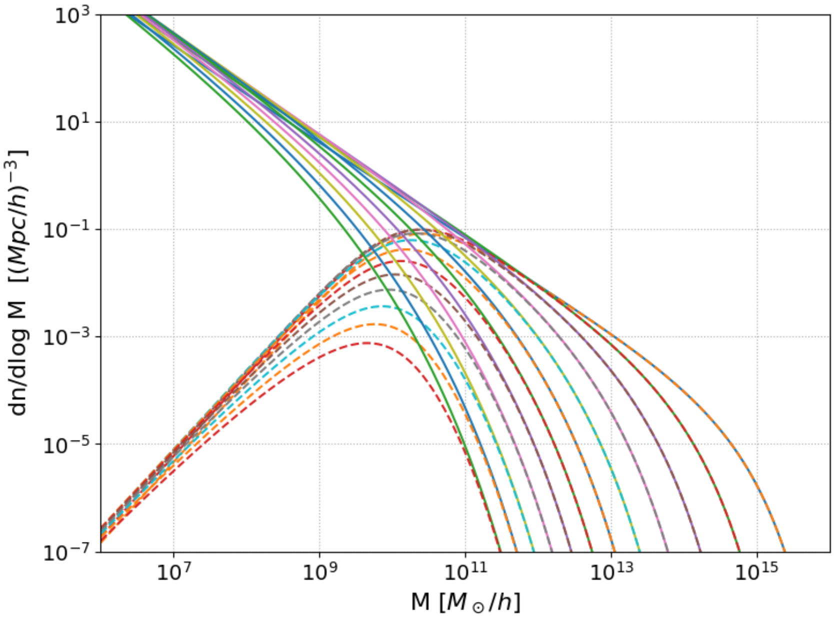

FIG. 1: Halo mass functions (HMFs) for CDM (solid lines) and ULDM (dashed lines, for $m_{22} = 1$) at redshifts $z = 0 \sim 11$, computed within the Press–Schechter formalism. HMFs increase as z decreases, with the highest curves corresponding to $z = 0$. Relative to CDM, ULDM shows a pronounced suppression of low-mass halos arising from the small-scale cutoff in the matter power spectrum.

The quantum Jeans mass scale is then defined by [11, 19, 28]

$$M_J(z) = \frac{4\pi}{3}\rho_{dm}(z)\,\lambda_J^3(z) = 1.204\times 10^8 \left(\frac{1+z}{m_{22}^2}\right)^{3/4} \left(\frac{\Omega_{\rm dm}}{0.27}\right)^{1/4} \left(\frac{h}{0.7}\right)^{1/2} M_\odot, \tag{II.3}$$

which implies that ULDM halos can provide sufficiently deep potential wells to confine baryons at high redshift. The dark matter halo mass must satisfy $M_h \geq M_J$, since the collapsing radius is constrained by $\lambda \geq \lambda_J$. The collapse is allowed only in the mass range $M_{\rm upper} \geq M_h \geq M_J$. The upper bound $M_{\rm upper}$ reflects the observed number density of galactic halos.

We shall assume $\Omega_{\rm dm} = 0.27$ and $h = 0.7$ in this paper. Then, the dark matter density at the redshift z is given by

$$\rho_{dm}(z) = 0.2485\times 10^{-29}\,(1+z)^3 {\rm g/cm}^3. \tag{II.4}$$

Taking $M_h \equiv \alpha^3\, M_J$, the collapse condition becomes $\alpha_{max}(z) \geq \alpha \geq 1$ where

$$\alpha_{max}(z) \equiv \left(\frac{M_{\rm upper}(z)}{M_J}\right)^{1/3}. \tag{II.5}$$

This leads to the halo mass as a function of the parameter α

$$M_h = \frac{4\pi}{3}\rho(z)\alpha^3\lambda_J^3(z) = 1.204\times 10^8\,\alpha^3 \left(\frac{1+z}{m_{22}^2}\right)^{3/4} M_\odot. \tag{II.6}$$

From the halo mass, to obtain M_{sol}, one may use the empirical formula for the core-halo relation for ULDM [29]

$$M_{sol} = 1.680\times 10^9 m_{22}^{-1}\,\frac{\xi'(z)}{\xi'(7)} \left(\frac{M_h}{10^{11}\,M_\odot}\right)^{1/3} M_\odot, \tag{II.7}$$

where $\xi'(z) = \sqrt{1+z}\,\xi(z)^{1/6}$ and $\xi(z) = (18\pi^2 + 82(\Omega_m(z)-1) - 39(\Omega_m(z)-1)^2)/\Omega_m(z) \simeq 177$ for large z. When $z > 7$ (or even $z > 1$), the relative change in $\xi(z)$ can be ignored. Taking the halo mass as given above, one has the core soliton mass

$$M_{sol} = 0.6319 \times 10^8\,\alpha \left(\frac{1+z}{m_{22}^2}\right)^{3/4} M_\odot. \tag{II.8}$$

The corresponding half-mass radius of the soliton is given by

$$r_{1/2} = \frac{2\beta f_0}{3}\frac{\hbar^2}{m^2}\frac{1}{GM_{sol}} = 3.539\,\frac{\beta}{m_{22}^2\alpha}\left(\frac{1+z}{m_{22}^2}\right)^{-3/4} \text{kpc} \tag{II.9}$$

with $f_0 = 3.9251$ [24]. We insert here an extra factor $2\beta/3$ to account for the reduction of the half-mass radius due to the presence of the baryonic matter. This factor β may be estimated to be around unity and below we shall set it to be unity for simplicity. To reproduce the observed galactic core sizes, we require $r_{1/2} \simeq$ kpc, which corresponds to $m \simeq 10^{-22}$eV [30]. Interestingly, the average dark matter density at the half mass radius

$$\rho_{1/2} = \frac{3M_{sol}}{8\pi r_{1/2}^3} = 1.152 \times 10^{-26}\alpha^4(1+z)^3\text{g/cm}^3 = 1.702 \times 10^{-4}\alpha^4(1+z)^3\text{M}_\odot/\text{pc}^3 \tag{II.10}$$

becomes independent of m. One finds that compared to the CDM model ULDM halos at redshift z can develop extremely compact and high-density cores, particularly in systems hosting massive solitons with large M_{sol}.

In this study, the halo mass function (HMF) is required to link the theoretically allowed collapse window to the cosmological abundance of host halos, thereby enabling a direct assessment of whether the proposed seed-formation conditions can account for the observed population of SMBHs. We adopt the Press–Schechter (PS) formalism [31], which gives

$$\frac{dn}{d\ln M} = -\frac{1}{2}\frac{\rho_0}{M}f(\nu)\frac{d\ln\sigma^2}{d\ln M}, \tag{II.11}$$

where $\nu \equiv \frac{\delta_c}{\sigma}$ with the critical overdensity δ_c, and dn is the abundance of halos within a mass interval dM. For the function $f(\nu)$ we use the model of Sheth & Tormen [32] for simplicity:

$$f(\nu) = A\sqrt{\frac{2}{\pi}}\sqrt{q\nu}\left(1 + (q\nu)^{-2p}\right)\exp\left[-\frac{q\nu^2}{2}\right] \tag{II.12}$$

with parameters $\{A = 0.3222, p = 0.3, q = 0.707\}$. Compared to the CDM case, the HMF for ULDM is suppressed at low masses as a result of the quantum pressure-induced cutoff in the matter power spectrum as [33]

$$\left.\frac{dn}{d\ln M}\right|_{\text{ULDM}} = \left.\frac{dn}{d\ln M}\right|_{\text{CDM}}\left[1 + \left(\frac{M_h}{M_0}\right)^{-1.1}\right]^{-2.2}, \tag{II.13}$$

where $M_0 = 1.6 \times 10^{10} m_{22}^{-4/3} M_\odot$. Fig. 1 shows HMFs for CDM and ULDM at various redshifts z, computed within the PS formalism [34]. Relative to CDM, ULDM shows a pronounced suppression of low-mass halos.

To determine M_{upper}, one should integrate the HMF and compare the resulting abundance with observations. At a given redshift z, we define the maximum halo mass $M_{upper}(z)$ as the mass above which the expected number of halos within the observable Universe over the full sky is exactly unity. To determine this quantity, following Ref. [35], we numerically compute the cumulative comoving number density

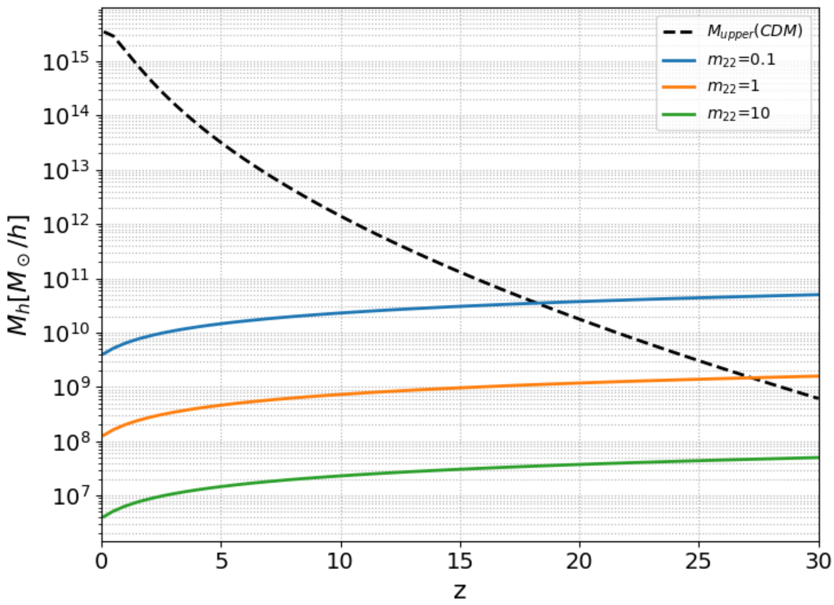

FIG. 2: The minimum mass M_J (the thick lines) and the maximum mass M_{upper} (the dashed line) of ULDM halos as functions of the redshift z for $m_{22} = (0.1, 1, 10)$, respectively. The halo mass decides the core masses of ULDM and those of gas clouds using the core-halo relation. For M_{upper}, we used the CDM approximation. At a given z, only halos with masses between M_J and M_{upper} can exist.

$n(> M, z)$ from the Sheth–Tormen halo mass function. We then multiply this by the comoving volume element corresponding to a thin redshift slice $dz = 0.5$ over the full sky, to obtain the expected number of halos $N(> M, z)$. The maximum halo mass is defined by solving the condition $N(> M_{upper}(z), z) = 1$. This construction represents the mass of the most massive halo that is most likely to be observed near that redshift. For simplicity, we use the CDM halo mass function rather than the ULDM one to compute M_{upper}, which is a good approximation in the large M_h regime. We also assume that the minimum halo mass at z is set by the Jeans mass $M_J(z)$. With these assumptions, the minimum and maximum halo masses at a given z can be estimated. Then, from M_J and $M_{\rm upper}$, one can obtain the ranges of α and M_{sol} using Eq. (II.7). Fig. 2 shows M_J and M_{upper} as functions of z for $m_{22} = (0.1, 1, 10)$, respectively. As z increases, M_{upper} decreases and the allowed range of M_h shrinks.

There is yet another upper bound for the soliton mass independent of z. The maximum gravitationally stable mass of an isolated ground-state configuration of ULDM is given by

$$M_{\rm max} \simeq 0.633\, \frac{M_{\rm Pl}^2}{m} \simeq \frac{8.5 \times 10^{11}}{m_{22}}\, M_\odot, \tag{II.14}$$

where $M_{\rm Pl}$ is the Planck mass. A soliton with a mass exceeding this scale is unstable to gravitational cooling and tends to lose mass [36], unless it is compressed under extreme conditions such as high-speed collisions or in the very high-redshift Universe. Consequently, we expect the maximum dark matter core mass to be effectively bounded at approximately M_{max}. This, in turn, may provide an upper bound on the masses of gas clouds within the core and on the ultimate masses of black hole formed and grown in the gas clouds. Note that M_{max} is usually different from the soliton mass in halos with mass M_{upper}.

III. BARYONIC COLLAPSE AND SEED-MASS SCALE INSIDE ULDM SOLITONS

Let us now place baryons inside the dark matter soliton. Realistically collapsing atomic gas generically develops a finite-density isothermal core owing to shocks and efficient radiative cooling. Therefore, a pseudo-isothermal profile

$$\rho_b(r) = \frac{\rho_0}{1 + r^2/r_c^2} \tag{III.1}$$

is the physically relevant description for the gas cloud [37], where ρ_0 is the central density and r_c is the core radius of the gas cloud. For the peudo-isothermal profile adopted in this work, the baryon mass up to radius r is given by

$$M_b(r) = 4\pi\rho_0\, r_c^3 \left(\frac{r}{r_c} - \arctan\frac{r}{r_c}\right) = 4\pi\rho_0\, r_c^3 G\left(\frac{r}{r_c}\right), \tag{III.2}$$

where $G(x) = x - \arctan(x)$. We adopt this mass profile for initial gas distribution within the ULDM cores. The concentration for the pseudo-isothermal profile $c_r \equiv r_{vir}/r_c$ can be about 30 [38], where r_{vir} is the virial radius of the gas clouds. Since we expect baryons to collapse into the gravitational potential well of ULDM cores, for simplicity, we assume r_{vir} is equal to the half-mass radius $r_{1/2}$ of the ULDM soliton cores. Then, $M_b(r_{vir}) \simeq 0.2M_{sol}/2 \simeq 4\pi\rho_0 r_c^3 G(r_{vir}/r_c)$, where 0.2 represents the approximate mass ratio of baryon to dark matter.

From Eq. (II.9), the core radius of baryon becomes

$$r_c = r_{1/2}/c_r \simeq 118.0 \times \frac{1}{\alpha m_{22}^2}\left(\frac{1+z}{m_{22}^2}\right)^{-3/4} \text{pc} \tag{III.3}$$

for $c_r = 30$. Then, one can obtain

$$M_b(r) \simeq \frac{0.1 M_{sol} G\left(\frac{r}{r_c}\right)}{G\left(c_r\right)} \tag{III.4}$$

and the baryon mass within the core

$$M_b(r_c) \simeq 4.764 \times 10^4 \alpha \left(\frac{1+z}{m_{22}^2}\right)^{3/4} M_\odot. \tag{III.5}$$

The average baryon density at the radius r_c becomes

$$\rho_b(r_c) = \frac{3M_b(r_c)}{4\pi r_c^3} \simeq 6.929 \times 10^{-3}\alpha^4(1+z)^3\, \mathrm{M_\odot/pc^3}. \tag{III.6}$$

We shall take this gas cloud with $r < r_c$ as a baryonic core that triggers collapse via the Jeans instability. The ULDM density at $r = r_c$ is much lower than the baryon density and its mass and gravitational contribution may be ignored once the gas core is formed.

The Jeans scale of ULDM determines the minimum halo mass able to collapse at high redshift, while the empirical soliton–halo relation fixes the soliton mass and dark matter core radius. These, in turn, set the baryonic core density, inflow velocity, shock temperature, and accretion rate.

The seed black hole mass may be expressed as

$$M_{bh} = f_{bh} M_b(r_c), \tag{III.7}$$

where we adopt a value $f_{bh} \simeq O(0.1)$, motivated by the fact that M_{bh} can grow to as much as about half of the quasi-star mass [39, 40]. Therefore, the mass of the black hole seeds formed at z satisfies

$$4.764 f_{bh} \times 10^4 \left(\frac{1+z}{m_{22}^2}\right)^{3/4} M_\odot \lesssim M_{bh} \lesssim 4.764 f_{bh} \times 10^4 \alpha_{max}\left(\frac{1+z}{m_{22}^2}\right)^{3/4} M_\odot, \tag{III.8}$$

where α_{max} can be inferred from M_{upper}. However, we will show that not all seeds within this mass range can form if the temperature of the gas is not sufficiently high. (See Eq. III.18 below.)

One can also estimate the maximum mass of SMBHs grown from their seeds. If we assume that a seed forms and grows within the most massive physically allowed soliton (i.e., $M_{\rm sol} \simeq M_{\rm max}$) and consumes remaining baryonic mass enclosed within the core, then its mass can be as large as $0.2 f_{bh} M_{\rm max}$. This sets the characteristic mass scale of the most massive SMBHs (not just seeds),

$$M_{bh}^{\rm max} \equiv 0.2 f_{bh} M_{\rm max} \simeq \frac{1.7 \times 10^{11} f_{bh}}{m_{22}} M_\odot. \tag{III.9}$$

Interestingly, from the observational data $M_{bh}^{\rm max} \simeq 10^{10} M_\odot$ [2], one can again get $m \simeq 10^{-22} eV$ for $f_{bh} = 0.1$. Because such massive solitons and the corresponding SMBHs are rare, mergers among them are expected to be infrequent. Consequently, this mass scale is not expected to change significantly due to mergers in cosmic history (see Fig. 3).

We now estimate the shock temperature of the core gas cloud as

$$T_{\rm core-cloud} \simeq \frac{3\mu m_{\rm H}}{16 k_{\rm B}} v_{\rm in}^2 \simeq \frac{3\mu m_{\rm H}}{8 k_{\rm B}} \frac{G M_b(r_c)}{r_c}, \tag{III.10}$$

where v_{in} denotes the gas inflow velocity, m_H is the hydrogen mass, and we take $\mu = 1.22$ for definiteness. This leads to

$$T_{\rm core-cloud} \simeq 96.33\,(1+z)^{\frac{3}{2}} \frac{\alpha^2}{m_{22}} \,{\rm K}. \tag{III.11}$$

Radiation cooling down to the Jeans temperature of the core gas clouds is required to have a collapse. To estimate the Jeans temperature we proceed as follows. We assume that the gas clouds have an initial radius r_c with the mass $M_b(r_c)$ given in the above. We further assume that the clouds have a *uniform* density $\rho_b(r_c)$ for $r < r_c$. The Jeans temperature can be obtained from the virial condition

$$2E_{kinetic} + E_{potential} = 0 \tag{III.12}$$

which becomes

$$2 \cdot \frac{3 k_B T_J}{2} \frac{M_b(r_c)}{\mu m_H} - \frac{3}{5} \frac{G M_b^2(r_c)}{r_c} = 0 \tag{III.13}$$

where μm_H is the average mass of the particles. This leads to the Jeans temperature

$$T_J = \frac{4\pi G \mu m_H \, r_c^2 \rho_b(r_c)}{15\, k_B}. \tag{III.14}$$

With the above r_c and $\rho_b(r_c)$, one finds

$$T_J \simeq 51.38\,(1+z)^{\frac{3}{2}} \frac{\alpha^2}{m_{22}} \,{\rm K}. \tag{III.15}$$

For gravitational collapse to commence, the gas must first cool from $T_{\text{core-cloud}}$ down to this temperature threshold. Once sufficient cooling has occurred, collapse can proceed. If $T_J < 3 \times 10^4$K, the gas becomes prone to fragmentation, as discussed below. In contrast, if $T_J > 3 \times 10^4$K, fragmentation is suppressed and direct collapse becomes possible. This requires

$$\alpha^2 \frac{(1+z)^{3/2}}{m_{22}} \gtrsim 583.9 \tag{III.16}$$

or

$$\alpha \gtrsim \alpha_{th} = 24.16\frac{m_{22}^{1/2}}{(1+z)^{3/4}} \tag{III.17}$$

which leads to

$$M_b(r_c)|_{\alpha=\alpha_{th}} = 1.151 \times 10^6 m_{22}^{-1}\,\mathrm{M}_\odot. \tag{III.18}$$

This with the observational constraint for the minimum seed mass $10^4 M_\odot \lesssim M_{bh} \lesssim 10^5 M_\odot$ gives an interesting bound for the particle mass of ULDM

$$1.151 f_{bh} \times 10^{-21} eV \lesssim m \lesssim 1.151 f_{bh} \times 10^{-20} eV, \tag{III.19}$$

which is consistent with the ULDM particle mass $m \simeq 10^{-22} eV$ inferred from galactic dynamics for $f_{bh} \simeq 0.1$. For $m \gg 10^{-22} eV$, SMBH seeds are too small in mass, whereas for $m \ll 10^{-22} eV$ their masses are excessively large and the seeds form too late (See Fig. 2). Gas clouds with masses below the threshold given in Eq. (III.18) fail to form black hole seeds and instead are expected to form star clusters. Interestingly, ULDM with the fiducial mass $m \simeq 10^{-22} eV$ seems to explain the minimum mass and the maximum mass of observed SMBHs. (See Fig. 3.)

From Eq. (III.17) and Eq. (III.3), one finds that the baryonic core size is smaller than $r_c|_{\alpha=\alpha_{th}} = 4.884$ pc$/m_{22}$, and the gas radius is smaller than $r_{vir}|_{\alpha=\alpha_{th}} \simeq 146.0$ pc$/m_{22}$, which is consistent with LRD observations [41]. From Eq. (III.3) and Eq. (III.6), the mass inflow rate $\dot{M} = 4\pi r_c^2 \rho_b(r_c) v_{in}$ is given by

$$\dot{M} \simeq 2.310 \times 10^{-3} \alpha^3 m_{22}^{-3/2} (1+z)^{9/4}\,\mathrm{M}_\odot/\mathrm{yr} > 32.58\,\mathrm{M}_\odot/\mathrm{yr}, \tag{III.20}$$

where we used the fact $\alpha > \alpha_{th}$ for the last inequality. This strong inflow further ensures that the gas collapses monolithically to form a massive protostellar core, providing a dynamical pathway to SMBH seeds. Once SMBH seeds form, ULDM solitons can supply additional gravitational potential that enhances the Bondi accretion rate onto a growing black hole seed [25], thereby enhancing the growth of SMBHs. Note that we do not consider angular momentum transfer and radiation feedback here.

In our model, SMBH seeds can form at high $z > 10$, possibly alleviating the tension of the standard scenario of black hole formation. Assuming Eddington-limited accretion with a radiative efficiency of $\epsilon = 0.1$, the black hole mass evolves as [42]

$$M(z_f) = M_i \exp\left[\frac{t(z_f) - t(z_i)}{t_{\mathrm{Sal}}}\right], \tag{III.21}$$

where $t_{\mathrm{Sal}} \simeq 45\,\mathrm{Myr}$ is the Salpeter e-folding timescale. For example, for the quasar in the galaxy UHZ1 with mass $M(z_f) = 4 \times 10^7 M_\odot$ at $z_f = 10.1$ [43], starting from an initial seed mass $M_i = 10^5\,M_\odot$, the required growth factor is $M(z_f)/M_i = 400$. This implies $t(z_f) - t(z_i) \simeq 2.7 \times 10^8$yr, which corresponds to an initial redshift $z_i \simeq 19$. (See Fig. 3.) Therefore, a heavy seed of mass $10^5 M_\odot$ in our scenario can naturally grow to $\sim 4 \times 10^7 M_\odot$ by $z \simeq 10$, demonstrating that early SMBH growth is feasible without invoking super-Eddington accretion. Fig. 3 shows the mass spectrum of observed quasars and candidate SMBHs in LRDs, which is well reproduced by our theoretical predictions. Our model predicts that many LRDs with SMBHs having masses $M_{bh} \gtrsim 10^5 M_\odot$ exist even beyond $z > 7$.

One can also determine the Jeans radius for baryon gas in terms of the average density ρ_b, and the temperature of the clouds as

$$r_J = \left(\frac{15\,k_B T}{4\pi G \mu m_H \rho_b}\right)^{1/2}. \tag{III.22}$$

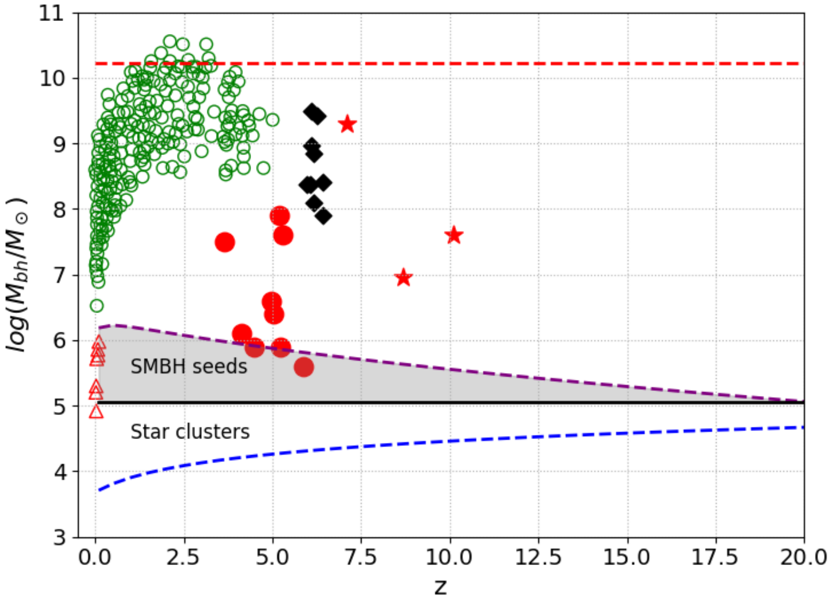

FIG. 3: The mass spectrum of SMBHs as a function of redshift. The black horizontal line denotes the masses of SMBH seeds corresponding to α_{th}. The blue dashed line and the purple dashed line denote the possible SMBH seed mass in Eq. (III.8) for $\alpha = 1$ and $\alpha = \alpha_{max}$, respectively. However, objects located between the black horizontal line and the blue line cannot be genuine black hole seeds because the gas temperature is too low and fragmentation is unavoidable; instead, these systems are expected to form star clusters. According to Eq. (III.17) in our model, only gas clouds corresponding to the SMBH seeds in the gray region possess sufficient self-gravity and temperature to form a SMBH seed. The dashed horizontal red line denotes the maximum black hole mass M_{bh}^{max} in Eq. (III.9) set by the stability of the soliton. Dots represent the observed masses of SMBHs from various observations [43–48]. The red triangles denote XMM-Newton observations [44], the green circles represent LBQS observations [45], the black diamonds denote CFHQS observations [46]. The red stars denote quasars with $z > 7$ [43, 47, 48], while the red star at $z = 10.1$ is for the quasar in the galaxy UHZ1. The red disks denote the candidate SMBHs in LRDs [9]. We choose $m = 10^{-22}eV$ and $f_{bh} = 0.1$ here.

Hence the baryon Jeans mass for a given density and temperature becomes

$$M_{bJ} = \frac{4\pi}{3}\rho_b r_J^3 = \left(\frac{375}{4\pi}\right)^{1/2}\left(\frac{k_B T}{G\mu m_H}\right)^{3/2}\rho_b^{-\frac{1}{2}}. \tag{III.23}$$

During the collapse, we assume the adiabatic process. Then one has the temperature

$$T = T_i\left(\frac{\rho_b}{\rho_i}\right)^{\gamma-1} \tag{III.24}$$

as before where T_i and ρ_i denote the initial temperature and density of the collapse. As the collapse goes on, the density becomes larger and then temperature increases since $\gamma > 1$ in general. Using this, one finds the updated baryon Jeans mass becomes

$$M_{bJ} = M_i\left(\frac{\rho_b}{\rho_i}\right)^{\frac{3}{2}(\gamma-4/3)}. \tag{III.25}$$

Therefore, if $\gamma < 4/3$, $M_{bJ} < M_i = M_b(r_c)$ and the fragmentation occurs as the collapse goes on. On the other hand, if $\gamma > 4/3$, $M_{bJ} > M_i = M_b(r_c)$ and the fragmentation is not allowed. When $T > 3\times10^4\,$K, $\gamma \simeq 5/3$ and there is no fragmentation. If $T < 3\times10^4\,$K, eventually the temperature passes through the range $T \sim 10^4$–3×10^4 K where $\gamma \simeq 1.1 \sim 1.3$ [3]. As a result, fragmentation into smaller stars occurs. When the collapse continues without fragmentation, a massive star with a mass on the order of $10^5 M_\odot$ can form. This star can subsequently collapse into a black hole, as its lifetime is very short.

IV. DISCUSSION

Observations of LRDs indicate that black hole seeds with masses $M_{\rm bh} \gtrsim 10^5 M_\odot$ already exist at $z \gtrsim 5$, and our scenario provides a physical interpretation of their origin. In this work, we show that this characteristic mass scale arises naturally in the ULDM scenario: when the baryonic core inside a soliton satisfies the temperature threshold, the enclosed gas mass reaches $M_b(r_c)\big|_{\alpha=\alpha_{\rm th}} \sim 10^6 M_\odot$, and a fraction $f_{\rm bh} \sim 0.1$ of this mass is converted into a black hole, yielding a seed of order $10^5 M_\odot$ without fine-tuning. A key outcome of this work is that the typical seed mass is not an ad hoc assumption but instead follows from the intrinsic mass and length scales set by ULDM.

The deep gravitational potential of the soliton also leads to rapid gas inflow and shock heating up to $T_{\text{core-cloud}} \gtrsim 3 \times 10^4$ K, suppressing molecular hydrogen cooling and preventing fragmentation, consistent with the compact, hot, and ionized environments inferred for LRDs. Below the threshold mass, gas clouds generically fail to form black hole seeds and are instead expected to collapse into dense star clusters, possibly accounting for the observed coexistence of AGN-like and young stellar-dominated LRDs [49], and potentially triggering the early formation of galaxies within a single physical framework. Furthermore, because the relevant Jeans scale and the soliton–halo relation favor the formation of dense cores at high redshift, our scenario predicts that SMBH seed formation is efficient at $z > 5$, matching the redshift range where LRDs are observed by JWST. Finally, we demonstrate that standard Eddington-limited accretion is sufficient to grow these seeds to the observed masses by $z \sim 10$, without invoking super-Eddington accretion.

This mechanism has several important implications for early structure formation. It provides a new connection between the small-scale physics of ULDM and the early growth of supermassive black holes, suggesting that the existence and typical masses of high-redshift SMBHs may carry indirect information about ULDM particle properties. At the same time, the scenario makes testable predictions for the host halo masses, event rates, and thermodynamic conditions of early black hole formation, which can be confronted with upcoming observations of the high-redshift Universe. For example, our model predicts a substantial population of LRDs hosting SMBHs with masses $M_{bh} \gtrsim 10^5 M_\odot$ at redshifts $z > 7$.

The effects of angular momentum and radiative feedback are not included in the present semi-analytic treatment, although the solitonic potential is expected to promote rapid infall. While our model qualitatively reproduces key properties inferred for LRDs and SMBHs, including the characteristic mass scale, formation epoch, and the presence of hot, ionized gas, a more quantitative comparison with their number density and host-halo properties is still needed. Addressing these issues will require high-resolution simulations and improved observational constraints, and the present work should therefore be regarded as a starting point for more detailed studies linking the small-scale physics of ULDM to the observed population of high-redshift SMBHs.

Acknowledgments

DB was supported in part by Basic Science Research Program through NRF funded by the Ministry of Education (2018R1A6A1A06024977).

---

[1] M. Volonteri, The Astronomy and Astrophysics Review **18**, 279 (2010), 1003.4404.
[2] K. Inayoshi, E. Visbal, and Z. Haiman, Annual Review of Astronomy and Astrophysics **58**, 27 (2020), 1911.05791.
[3] M. A. Latif and A. Ferrara, Publications of the Astronomical Society of Australia **33**, e051 (2016), 1605.07391.
[4] V. Bromm and A. Loeb, The Astrophysical Journal **596**, 34 (2003), astro-ph/0212400.
[5] M. C. Begelman, M. Volonteri, and M. J. Rees, Monthly Notices of the Royal Astronomical Society **370**, 289 (2006), astro-ph/0602363.
[6] J. A. Regan and M. G. Haehnelt, Monthly Notices of the Royal Astronomical Society **396**, 343 (2009), 0810.2802.
[7] C. Shang, G. L. Bryan, and Z. Haiman, Monthly Notices of the Royal Astronomical Society **402**, 1249 (2010), 0906.4773.
[8] E. Visbal, Z. Haiman, and G. L. Bryan, Monthly Notices of the Royal Astronomical Society **445**, 1056 (2014).
[9] V. Rusakov, D. Watson, G. P. Nikopoulos, et al., Nature **649**, 574 (2026).
[10] L. E. Strigari, J. S. Bullock, M. Kaplinghat, J. D. Simon, M. Geha, B. Willman, and M. G. Walker, Nature **454**, 1096 (2008), 0808.3772.
[11] J.-W. Lee, Phys. Lett. B **756**, 166 (2016), 1511.06611.
[12] M. R. Baldeschi, G. B. Gelmini, and R. Ruffini, Physics Letters B **122**, 221 (1983).
[13] S.-J. Sin, Phys. Rev. D **50**, 3650 (1994), hep-ph/9205208.
[14] J.-W. Lee and I.-G. Koh, Phys. Rev. D **53**, 2236 (1996), hep-ph/9507385.
[15] T. Matos and F. S. Guzman, Class. Quant. Grav. **17**, L9 (2000), gr-qc/9810028.
[16] W. Hu, R. Barkana, and A. Gruzinov, Phys. Rev. Lett. **85**, 1158 (2000), astro-ph/0003365.
[17] C. G. Boehmer and T. Harko, Journal of Cosmology and Astroparticle Physics **2007**, 025??25 (2007).
[18] P. H. Chavanis, Astronomy and Astrophysics **537**, A127 (2012).
[19] L. Hui, J. P. Ostriker, S. Tremaine, and E. Witten, Phys. Rev. D **95**, 043541 (2017), 1610.08297.
[20] T. Matos, L. A. Ure참a-L처pez, and J.-W. Lee, Frontiers in Astronomy and Space Sciences **11** (2024).
[21] E. G. M. Ferreira, The Astronomy and Astrophysics Review **29**, 7 (2021), 2005.03254.
[22] D. J. E. Marsh, Phys. Rep. **643**, 1 (2016), 1510.07633.
[23] H.-Y. Schive, T. Chiueh, and T. Broadhurst, Nature Physics **10**, 496 (2014), 1406.6586.
[24] L. Hui, J. P. Ostriker, S. Tremaine, and E. Witten, Physical Review D **95**, 043541 (2017), 1610.08297.
[25] H.-H. S. Chiu, H.-Y. Schive, H.-Y. K. Yang, H. Huang, and M. Gaspari, Phys. Rev. Lett. **134**, 051402 (2025).
[26] H. Koo, D. Bak, I. Park, S. E. Hong, and J.-W. Lee, Physics Letters B **856**, 138908 (2024).
[27] K. Omukai, The Astrophysical Journal **546**, 635 (2001), astro-ph/0011446.
[28] M. Y. Khlopov, B. A. Malomed, I. B. Zeldovich, and Y. B. Zeldovich, Mon. Not. Roy. Astron. Soc. **215**, 575 (1985).
[29] H. H. S. Chiu, H.-Y. Schive, H.-Y. K. Yang, H. Huang, and M. Gaspari, Phys. Rev. Lett. **134**, 051402 (2025), 2501.09098.
[30] H.-Y. Schive, T. Chiueh, and T. Broadhurst, Nature Phys. **10**, 496 (2014), 1406.6586.
[31] W. H. Press and P. Schechter, Astrophys. J. **187**, 425 (1974).
[32] R. K. Sheth, H. J. Mo, and G. Tormen, Mon. Not. Roy. Astron. Soc. **323**, 1 (2001), astro-ph/9907024.
[33] H.-Y. Schive, T. Chiueh, T. Broadhurst, and K.-W. Huang, The Astrophysical Journal **818**, 89 (2016).
[34] S. Murray, C. Power, and A. Robotham (2013), 1306.6721.
[35] I. Harrison and P. Coles, Mon. Not. Roy. Astron. Soc. **421**, L19 (2012), 1111.1184.

[36] S. L. Liebling and C. Palenzuela, Living Rev. Rel. **26**, 1 (2023), 1202.5809.
[37] A. Trinca, A. Lupi, Z. Haiman, M. Volonteri, R. Valiante, R. Schneider, and R. Decarli (2026), 2601.14370.
[38] P. R. Shapiro, I. T. Iliev, and A. C. Raga, Monthly Notices of the Royal Astronomical Society **307**, 203??24 (1999).
[39] J. Hassan, R. Perna, M. Cantiello, P. Armitage, M. Begelman, and T. Ryu (2025), 2510.18301.
[40] A. D. Santarelli, C. B. Campbell, E. Farag, E. P. Bellinger, P. Natarajan, and M. E. Caplan (2025), 2510.11772.
[41] C. A. Guia, F. Pacucci, and D. D. Kocevski, Research Notes of the AAS **8**, 207 (2024), ISSN 2515-5172.
[42] J. W. Moffat (2020), 2011.13440.
[43] P. Natarajan, F. Pacucci, A. Ricarte, A. Bogdan, A. D. Goulding, and N. Cappelluti (2023), 2308.02654.
[44] G. C. Dewangan, S. Mathur, R. E. Griffiths, and A. R. Rao, The Astrophysical Journal **689**, 762 (2008).
[45] M. Vestergaard and P. S. Osmer, The Astrophysical Journal **699**, 800??16 (2009).
[46] C. J. Willott, L. Albert, D. Arzoumanian, J. Bergeron, D. Crampton, P. Delorme, J. B. Hutchings, A. Omont, C. Reyl쳅, and D. Schade, The Astronomical Journal **140**, 546??60 (2010).
[47] D. J. Mortlock, S. J. Warren, B. P. Venemans, M. Patel, P. C. Hewett, R. G. McMahon, C. Simpson, T. Theuns, E. A. Gonz챠les-Solares, A. Adamson, et al., Nature **474**, 616??19 (2011).
[48] R. L. Larson, S. L. Finkelstein, D. D. Kocevski, T. A. Hutchison, J. R. Trump, P. Arrabal Haro, V. Bromm, N. J. Cleri, M. Dickinson, S. Fujimoto, et al., The Astrophysical Journal Letters **953**, L29 (2023).
[49] M. Carranza-Escudero, C. J. Conselice, N. Adams, T. Harvey, D. Austin, P. Behroozi, L. Ferreira, K. Ormerod, Q. Duan, J. Trussler, et al., The Astrophysical Journal Letters **989**, L50 (2025).