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Boosting Supermassive Black Hole Growth in the Early Universe by Fuzzy Dark Matter Solitons

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Observations of massive supermassive black holes (SMBHs) in the early universe challenge existing black hole formation models. We propose that soliton cores in fuzzy dark matter (FDM) offer a potential solution to this timing problem. Our FDM cosmological zoom-in simulations confirm that for a particle mass $m_{\text{FDM}} \sim 10^{-22}$ eV, solitons are well developed at redshift $z \sim 7$ with masses of $\sim 10^9 M_\odot$, comparable to the observed SMBHs. We then demonstrate using hydrodynamic simulations that, compared to cold dark matter, these high- z massive FDM solitons with mass M_s can provide additional gravitational potential to accrete gas and boost the Bondi accretion rate of a growing black hole seed with mass M_{BH} by up to two to four orders of magnitude, in the regime of efficient cooling and negligible radiation pressure. This accretion boosting mechanism is effective for 10^{-22} eV $\lesssim m_{\text{FDM}} \lesssim 10^{-20}$ eV and potentially beyond as long as $M_s > M_{\text{BH}}$.

Introduction.—Observations of supermassive black holes (SMBHs) at high redshifts $z \gtrsim 7$ [1–8] with masses of $M_{\text{BH}} \gtrsim 10^7\text{--}10^9 M_\odot$ represent a significant challenge to existing black hole (BH) formation models [9–12]. These SMBHs may originate from various BH seeds, including massive seeds exceeding $10^4 M_\odot$ formed through direct gas collapse [13–15], lighter seeds ($\sim 10^2\text{--}10^3 M_\odot$) formed from Population III stars [16–18], or intermediate-mass seeds ($\sim 10^3\text{--}10^4 M_\odot$) formed from runaway collisions in dense star clusters [19, 20]. To reach the observed SMBH mass at such high redshifts, cosmological simulations indicate that these BH seeds must undergo rapid accretion at rates exceeding those predicted by standard accretion models. The accretion of heavy seeds must exceed the Bondi accretion rate [21, 22], which characterizes the steady-state, spherically-symmetric, and adiabatic accretion onto a point mass from a uniform gas [23][112]. In contrast, lighter seeds need to accrete at or above the Eddington limit [11, 24–29] for extended periods. Maintaining such high accretion rates is challenging because the Bondi accretion rate can only exceed the Eddington limit in extremely dense environments [30], and feedback from the active galactic nucleus (AGN) further alters the mass accretion [31–34].

Ultra-light dark matter (ULDM) [35–43], also known as scalar field dark matter or wave dark matter, is one of the leading dark matter candidates. It consists of ultra-light bosons with a broad particle mass range of $\sim 10^{-22}\text{--}10^{-6}$ eV and is well motivated by quantum chromodynamics axions [44] and string theory [45]. Extensive

studies have examined ULDM through various astrophysical probes [46–65] and ground-based experiments [66–74]. The high sensitivity of gravitational-wave detectors also offers a promising detection method [75–79]. Furthermore, heating from the decay of axion-like particles may inhibit molecular hydrogen formation, potentially leading to the direct collapse of gas clouds into SMBHs [80].

ULDM in the mass range of $m_{\text{FDM}} \sim 10^{-22}\text{--}10^{-20}$ eV is also referred to as fuzzy dark matter (FDM) [81], characterized by distinctive wave-like features on sub-kiloparsec scales, such as density granulation and the suppression of low-mass halos [82]. It may address small-scale challenges associated with cold dark matter (CDM) [83, 84], which remain a topic of debate (e.g., [85]). In particular, while maintaining the same large-scale structure as CDM, FDM predicts a unique structure at the center of each dark matter halo — a soliton core, which is a stationary, ground-state solution of the Schrödinger-Poisson equation [86–88]. The soliton half-density radius r_s and total mass M_s as a function of the host halo mass M_h and redshift z , derived from both numerical simulations [86, 87, 89, 90] and analytical models assuming thermal equilibrium between halos and solitons [87, 91, 92], are

$$r_s = 0.135 m_{22}^{-1} \left(\frac{\zeta'(z)}{\zeta'(7)} \right)^{-1} \left(\frac{M_h}{10^{11} M_\odot} \right)^{-1/3} \text{ kpc} \quad (1)$$

and

$$M_s = 1.68 \times 10^9 m_{22}^{-1} \left(\frac{\zeta'(z)}{\zeta'(7)} \right) \left(\frac{M_h}{10^{11} M_\odot} \right)^{1/3} M_\odot, \quad (2)$$

where $m_{22} \equiv m_{\text{FDM}}/10^{-22} \text{ eV}$, $\zeta' = (1+z)^{1/2} \zeta(z)^{1/6}$ with $\zeta(z) = (18\pi^2 + 82(\Omega_m(z) - 1) - 39(\Omega_m(z) - 1)^2)/\Omega_m(z) \sim 180$ at $z \geq 1$ [93], and Ω_m is the matter density parameter. M_s is roughly 4.2 times larger than the mass enclosed within r_s . Importantly, it remains uncertain whether Eqs. (1) and (2) can be reliably applied to massive halos at high redshifts, necessitating the FDM cosmological simulations performed in this study.

In this Letter, we first conduct FDM cosmological zoom-in simulations with $m_{22} \sim 1$ to confirm that FDM halos with $M_h \sim 10^{11} M_\odot$ at $z \sim 7$ have well-developed solitons with masses of $\sim 10^9 M_\odot$, comparable to SMBHs and in line with Eqs. (1) and (2). We then demonstrate using hydrodynamic simulations that, compared to CDM, these high- z massive FDM solitons can provide additional gravitational potential to accrete gas and boost the BH mass accretion rate $\dot{M}_{\text{BH}}$ by up to two to four orders of magnitude when cooling is efficient and radiation pressure is ignored. This additional boost provided by FDM solitons offers a potential explanation for the rapid growth of SMBHs in the early universe, with accretion rates significantly exceeding those predicted by the conventional Bondi model.

High- z FDM solitons.—We conduct FDM cosmological zoom-in simulations targeting a halo with $M_h \sim 7 \times 10^{10} M_\odot$ at $z = 6.87$ using the GPU-accelerated adaptive refinement code **GAMER** [94]. We investigate $m_{22} = 1.6$ and 0.8 in a cubic domain with a comoving side length of 15 Mpc, a 512^3 root grid, and up to nine refinement levels, achieving a maximum comoving resolution of 57 pc. Both the simulation domain length and maximum resolution are about an order of magnitude greater than those in [88]. The code employs a new hybrid scheme integrating a fluid formulation on coarser grids for large scales and a wave formulation on refined grids for small scales [95] (see also [96]). We construct the initial conditions at $z = 100$ using **AxionCAMB** [97] and **MUSIC** [98]. For comparison, we also conduct CDM simulations using the code **GADGET-2** [99] with 1024^3 particles and with a similar initial condition, except that there is no small-scale suppression by FDM quantum pressure in the initial power spectrum.

Fig. 1 shows the density and gravitational potential distributions of this target halo. FDM halos with $m_{22} = 0.8$ (1.6) formed at $z \sim 9$ (11) and had well-developed solitons by $z \sim 7$ with $r_s \sim 0.25$ (0.10) kpc and $M_s \sim 1.4 \times 10^9$ (8.7×10^8) $M_\odot$, consistent with Eqs. (1) and (2). The larger the m_{22} , the earlier the halos and solitons form. These solitons match well the analytical soliton profile [86] and result in a significantly deeper central gravitational potential than the CDM halo. Additionally, our target halos, with $M_h \sim 7 \times 10^{10} M_\odot$

at $z \sim 7$, are common in both CDM and FDM with $m_{22} \gtrsim 1$, as this mass scale exceeds the FDM half-mode mass, $M_{1/2} = 3.8 \times 10^{10} m_{22}^{-4/3} M_\odot$, below which the FDM halo mass function drops significantly compared to CDM [100]. These validations reinforce the soliton setup in the following hydrodynamic simulations for investigating how this additional gravitational potential of solitons may boost gas accretion and SMBH growth in the early universe.

BH-soliton accretion setup.—After confirming the existence of high- z solitons, we examine how the soliton gravity boosts the accretion rate of a growing BH seed. We perform hydrodynamic simulations with **GAMER** to investigate mass accretion from a uniform gas initially at rest onto a central BH embedded in an external soliton potential ϕ_{s0} [89]. To measure $\dot{M}_{\text{BH}}$, we adopt the well-tested prescription in [101], including a spherical ‘void’ region to model the BH horizon, within which we reset the density, temperature, and velocity to negligible values after each time-step Δt , and infer $\dot{M}_{\text{BH}}$ from the mass difference over Δt . In order to find the steady-state solution, we fix M_{BH} and perform simulations until $\dot{M}_{\text{BH}}$ saturates (~ 1000 Bondi times). The gas self-gravity is ignored as it is much weaker than the BH and soliton gravity.

To minimize numerical errors associated with the rapid accretion caused by a massive soliton, we gradually introduce the soliton potential by $\phi_s(r, t) = \phi_{s0}(r) \tanh(t/\beta t_g)$, where $t_g = GM_s/4.2c_s^3$, G is the gravitational constant, and c_s is the ambient sound speed. $\beta \sim 0.1$ – 10 is a dimensionless parameter to ensure the mass accretion remains adiabatic during the entire simulation.

Table I lists the adopted simulation parameters. We use an ideal gas equation of state with a range of effective adiabatic index $\gamma = 1$ – $5/3$, mimicking the composite effect of non-adiabatic micro-scale processes, which are often dominated by radiative cooling in astrophysical halos. We probe the ambient temperature $T_0 = 1$ – $10 \text{ eV} \sim 10^4$ – 10^5 K to account for different primordial environments [102–105]. To investigate how the soliton gravity enhances the accretion rate of a growing BH seed, we fix a typical seed of $M_{\text{BH}} = 10^5 M_\odot$ and probe $m_{22} = 1$ – 10 to ensure $M_s \gg M_{\text{BH}}$. Note that within this parameter space, $r_s \gg r_{\text{Bon}}$, indicating that the soliton and BH gravity dominate in the outer and inner regions, respectively.

We adopt a cubic computational domain with a side length of 8 kpc and zero-gradient boundary conditions. We choose a 64^3 root grid and double the resolution as the distance to the BH is halved, with up to 15 refinement levels, covering a dynamic range of six orders of magnitude with a maximum resolution of $3.8 \times 10^{-3} \text{ pc}$. This ensures that the void region for measuring $\dot{M}_{\text{BH}}$ is resolved by at least 12 cells. Simulations are conducted in three-dimensional Cartesian coordinates to facilitate

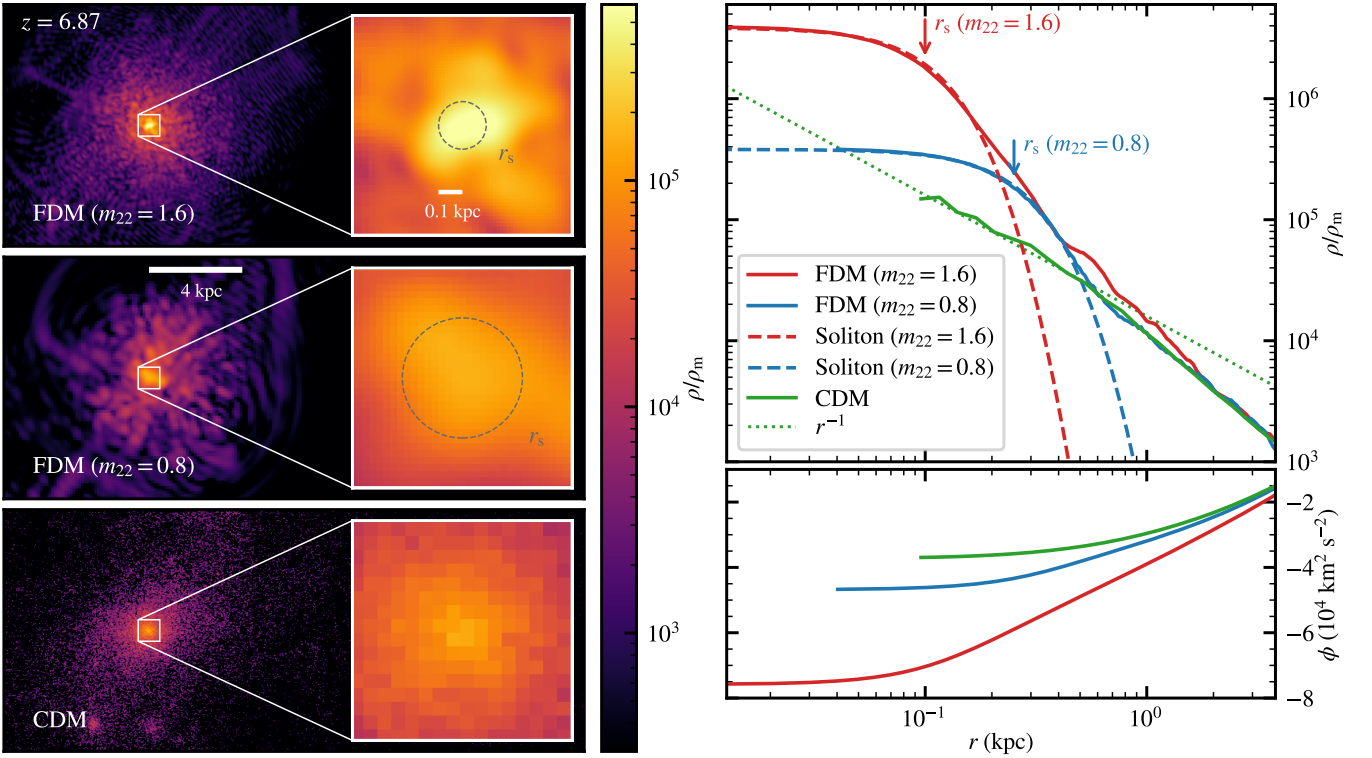


FIG. 1: FDM cosmological zoom-in simulations. (Left) Projected density in a 1 kpc thick slab through a $\sim 7 \times 10^{10} M_{\odot}$ halo with $m_{22} = 1.6$ (top) and 0.8 (middle) at $z = 6.87$, normalized to the mean matter density ρ_m ; a CDM counterpart is shown for comparison (bottom). Insets highlight the central regions, with dashed circles indicating the soliton radii. (Right) Density and gravitational potential profiles of the FDM and CDM halos (solid lines) and soliton solutions of [86] (dashed lines). Compared to CDM, distinct and massive soliton cores with $r_s \sim 0.1$ kpc and $M_s \sim 10^9 M_{\odot}$ form in FDM, in agreement with Eqs. (1) and (2), providing additional gravitational potential to accrete gas and boost SMBH growth in the early universe (see Figs. 2–4).

TABLE I: Simulation parameters. Parameters in the lower table are derived from those in the upper table, assuming $\gamma = 1.2$. Derived parameters are presented in ascending order of $T_0 = 1\text{--}10$ eV for Bondi-related parameters and $m_{22} = 1\text{--}10$ for soliton parameters.

Parameter	Value
BH mass (M_{BH})	$10^5 M_{\odot}$
Adiabatic index (γ)	$1\text{--}5/3$
Ambient density (ρ_0)	10^{-23} g/cm^3
Atomic mass fraction	0.62
Ambient temperature (T_0)	$1\text{--}10$ eV
FDM particle mass (m_{22})	$1\text{--}10$
Halo mass (M_h)	$10^{11} M_{\odot}$
Redshift (z)	7
Bondi radius (r_{Bon})	$2.3\text{--}0.23$ pc
Bondi time (t_{Bon})	$165\text{--}5.2$ kyr
Bondi sonic point	$0.81\text{--}0.081$ pc
Soliton radius (r_s)	$135\text{--}13.5$ pc
Soliton mass (M_s)	$16.8\text{--}1.68 \times 10^8 M_{\odot}$

future incorporation of non-spherically symmetric effects such as turbulence and rotation. We use the MUSCL-

Hancock integration scheme [106], the HLLC Riemann solver [106], and the piecewise parabolic reconstruction method [107]. Dual-energy formalism is applied to reduce truncation errors in supersonic flow.

BH-soliton accretion results.—Fig. 2 shows the gas profiles of a representative case with $\gamma = 1.2$, $T_0 = 10$ eV, and $m_{22} = 1$. We compare three scenarios: BH-only, soliton-only, and with both BH and soliton. The BH-only case confirms that our simulation setup can reproduce the analytical Bondi solution. The simulation with both BH and soliton shows a two-stage accretion attributed to $M_s \gg M_{\text{BH}}$ and $r_s \gg r_{\text{Bon}}$. The first-stage accretion results from the soliton gravity and is stalled around r_s due to its finite size. The gas density and temperature at $r_{\text{Bon}} \lesssim r \lesssim r_s$ increase by a factor of $\rho_{\text{inc}}/\rho_0 \sim 1000$ and $T_{\text{inc}}/T_0 \sim 4$, respectively, satisfying the adiabatic condition $T \propto \rho^{\gamma-1}$. Both ρ_{inc} and T_{inc} are independent of M_{BH} , as confirmed by the consistency between the soliton-only and BH-soliton simulations at $r \gtrsim r_{\text{Bon}}$.

The second-stage accretion occurs within r_{Bon} where the BH gravity dominates over the soliton gravity. The gas profiles still satisfy the Bondi solution but with a boosting factor $\alpha \sim 125$ compared to the BH-only case due to the significantly increased ambient density

from the first-stage accretion. Here we define $\alpha = \dot{M}_{\text{BH}}/\dot{M}_{\text{Bon0}}$, where $\dot{M}_{\text{BH}}$ is measured directly from a BH-soliton simulation and $\dot{M}_{\text{Bon0}}$ is the theoretical BH-only Bondi accretion rate from T_0 and ρ_0 . Note that from $\dot{M}_{\text{Bon}} \propto \rho T^{-3/2}$ and the adiabatic condition during the first-stage accretion, we can infer $\alpha = (\rho_{\text{inc}}/\rho_0)^{(5-3\gamma)/2} \sim 125$, in agreement with the measured value.

Fig. 3 plots the boosting factor α as a function of γ for the four corner cases: $(T_0/\text{eV}, m_{22}) = (1, 1), (1, 10), (10, 1), (10, 10)$, revealing a strong dependence on γ and a relatively weak dependence on T_0 and m_{22} . In general, the softer the equation of state, the larger the α , increasing from close to unity ($\gamma = 5/3$) to three to four orders of magnitude ($\gamma = 1$) regardless of T_0 and m_{22} . This is consistent with the adiabatic prediction, $\alpha = (\rho_{\text{inc}}/\rho_0)^{(5-3\gamma)/2}$, which gives $\alpha = 1$ for $\gamma = 5/3$ and $\alpha = \rho_{\text{inc}}/\rho_0$ for $\gamma = 1$. For comparison, we also plot the Eddington limit $\alpha = \dot{M}_{\text{Edd}}/\dot{M}_{\text{Bon0}}$ with $T_0 = 1$ and 10 eV, where $\dot{M}_{\text{Edd}}$ is the Eddington accretion rate assuming a radiation efficiency of 10%. Note that $\dot{M}_{\text{BH}} \gtrsim \dot{M}_{\text{Edd}}$ for $\gamma \lesssim 1.4$ at $T_0 = 1$ eV and $\gamma \lesssim 1.2$ at $T_0 = 10$ eV, indicating that solitons can help approach or even exceed the Eddington limit.

Fig. 3 suggests that α increases with increasing m_{22} and decreasing T_0 for $\gamma \lesssim 1.4$. To investigate this further, we plot in Fig. 4 the boosting factor as a function of T_0 and m_{22} for $\gamma = 1.2$, showing an approximately linear dependence on T_0^{-1} and m_{22} . This relation is attributed to the balance between gas pressure and soliton gravity in the first-stage accretion. For a given ρ_0 , a lower T_0 leads to higher density compression and thus a larger α . Similarly, for a fixed M_h in Eqs. (1) and (2), the gravitational acceleration near the soliton radius scales as $M_s/r_s^2 \propto (1+z)^{3/2}m_{22}$. Therefore, a larger m_{22} corresponds to a stronger soliton gravity that can accrete more gas inside r_s to boost $\dot{M}_{\text{BH}}$. The above $(1+z)^{3/2}$ dependence also indicates that for a given M_h , solitons have a greater impact at higher redshifts. This can compensate for the lower $\dot{M}_{\text{BH}}$ in the BH-only accretion scenario at higher redshifts when M_{BH} is low.

Discussion.—The two-stage accretion occurs only when $r_s \gg r_{\text{Bon}}$ and $M_s \gg M_{\text{BH}}$. The former condition can be violated for higher m_{22} or lower T_0 , assuming a fixed $M_{\text{BH}} = 10^5 M_\odot$. For example, for $m_{22} \sim 100$ and $T_0 = 1$ eV, $r_s \sim r_{\text{Bon}}$, and the accretion reverts to the single-stage process. However, $M_s \sim 10^7 M_\odot \gg M_{\text{BH}}$ remains ($M_s \propto m_{22}^{-1}$ from Eq. (2)), and our simulation demonstrates a boosting factor of ~ 1000 for $\gamma = 1.2$. For even larger particle masses in the range $10^3 \lesssim m_{22} \lesssim 10^4$, as suggested by some stringent constraints [47, 53, 108], substantial boosting is still expected for lighter BH seeds with $M_{\text{BH}} \lesssim 10^5 M_\odot$. As M_{BH} grows with time, eventually $M_{\text{BH}} \gtrsim M_s$, and $\dot{M}_{\text{BH}}$ transitions back to the conventional single-stage Bondi accretion without significant boosting.

The boosting factor α strongly depends on the effec-

tive adiabatic index γ , reaching $\alpha \gtrsim 10$ (100) only when $\gamma \lesssim 1.4$ (1.2) (Fig. 3). What is the appropriate value of γ in the environment of early BHs? Consider a gas halo with $T \sim 1\text{--}10$ eV, the normalized cooling rate is $\Lambda \sim 10^{-22}\text{--}10^{-21}$ erg cm³/s [107]. Assuming a gas number density of ~ 10 cm⁻³, the cooling timescale is $\sim 0.1\text{--}1$ kyr, significantly shorter than the Bondi time t_{Bon} in Table I. This estimate indicates efficient cooling and a low $\gamma < 5/3$ tied to a large boosting. This is further corroborated by the expectation of chaotic cold accretion frequently occurring in massive BHs and developing similarly elevated boosting factors [109].

To verify the robustness of our simulations, we confirm that the saturated α is insensitive to M_{BH} , ρ_0 , and initial gas profiles, provided that $M_{\text{BH}} \ll M_s$ holds and the gas self-gravity is negligible. Additionally, α is independent of the radius of the central void region as long as it is within the sonic point and resolved by at least 12 cells. The artificial parameter β , used for the gradual inclusion of the soliton potential, introduces $\sim 35\%$ uncertainty in α but does not affect the main conclusions of this study.

Rotation, turbulence, magnetic field, AGN feedback, and radiation pressure, not currently considered in this study, can provide additional support against accretion and reduce both BH-only and BH-soliton accretion rates. Nevertheless, the boosting factor, defined as the ratio between the two rates, will likely remain substantial. Our current three-dimensional simulations make it straightforward to incorporate these additional factors, along with radiative cooling, which we will consider in future work.

We note that the boosting mechanism proposed here is likely applicable to other dark matter models that can form dense cores at high redshifts, such as self-interacting dark matter [110].

Conclusions.—Soliton cores in FDM provide a potential solution to the timing problem of SMBH growth in the early universe. Our FDM cosmological zoom-in simulations with a particle mass $m_{22} \sim 1$ confirm that $\sim 10^{11} M_\odot$ halos at $z \sim 7$ have well-developed solitons with masses of $\sim 10^9 M_\odot$, comparable to the observed SMBH masses at similar epochs and significantly more massive than the BH seeds commonly employed in cosmological simulations. We then demonstrate quantitatively with hydrodynamic simulations how the additional soliton potential accretes gas to increase the ambient gas density around the BH, thereby boosting the Bondi accretion rate in its early evolution stage when the BH mass is much smaller than the soliton mass. We find that the BH accretion rate $\dot{M}_{\text{BH}}$ can be boosted by two to four orders of magnitude with $m_{22} = 1\text{--}100$, provided that radiative cooling is efficient leading to an effective $\gamma \lesssim 1.2$. Substantial boosting is also expected for $10^2 \lesssim m_{22} \lesssim 10^4$, especially for lighter BH seeds. For the same γ , $\dot{M}_{\text{BH}}$ increases approximately linearly with decreasing ambient gas temperature and increasing

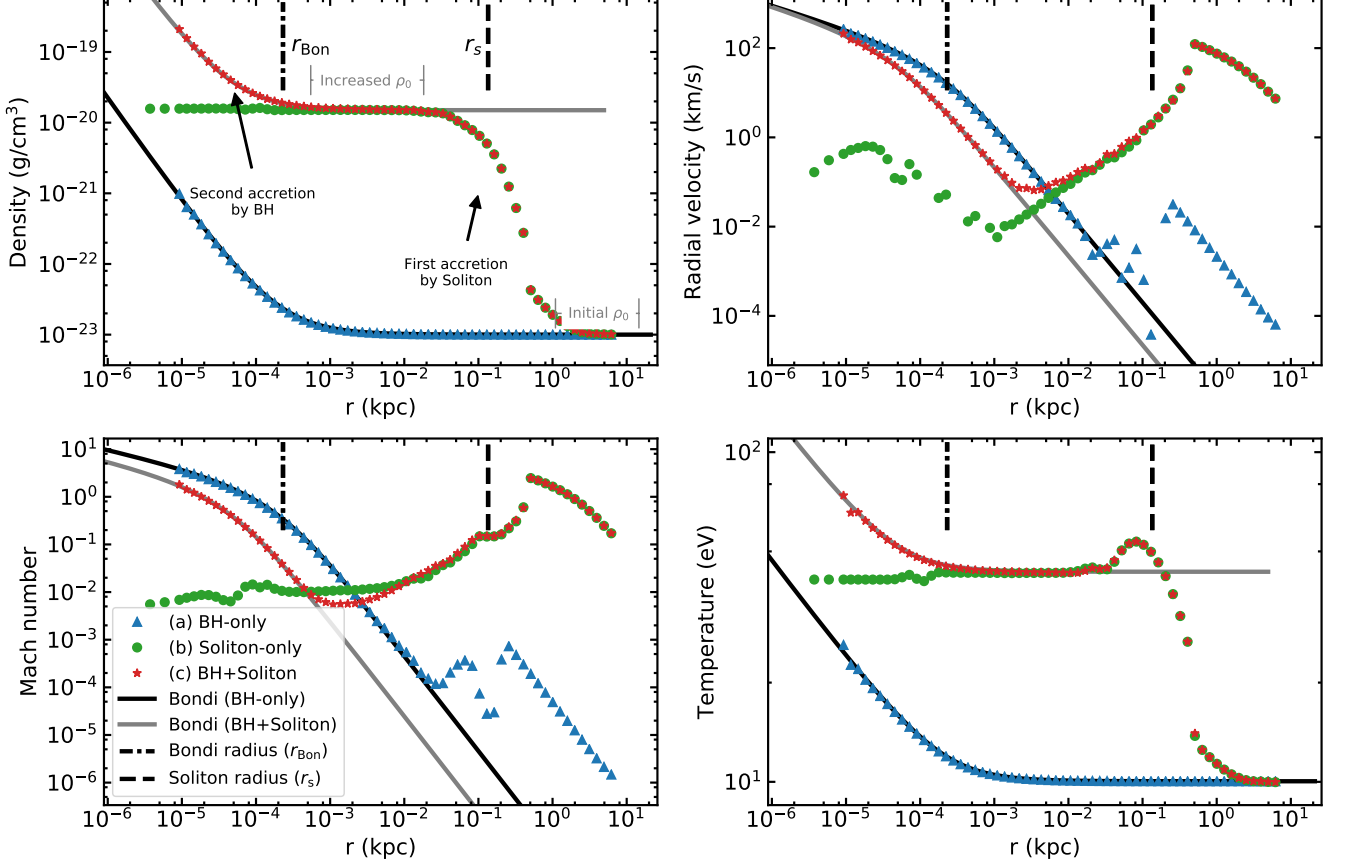


FIG. 2: Gas profiles of a representative BH-soliton accretion simulation with $\gamma = 1.2$, $T_0 = 10$ eV and $m_{22} = 1$. We compare three cases: (a) BH-only (b) soliton-only, and (c) both BH and soliton. Case (a) confirms that our simulation setup can reproduce the analytical Bondi solution (black solid line). Case (b) illustrates the region dominated by the soliton gravity. Case (c) presents the two-stage accretion feature of a BH-soliton system; the profile within r_{Bon} still matches the Bondi solution (grey solid line) but with a boosting factor $\alpha \sim 125$ (see Fig. 3) compared to the BH-only case due to the greatly increased ambient density in $r_{\text{Bon}} \lesssim r \lesssim r_s$ accreted by the soliton gravity. See text for details.

m_{22} . Such boosted accretion rates due to FDM could provide a natural way for early BHs to maintain near- or super-Eddington accretion and substantially aid their rapid growth in the early universe.

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Data availability.—Simulation source code is available at <https://github.com/gamer-project/gamer>.

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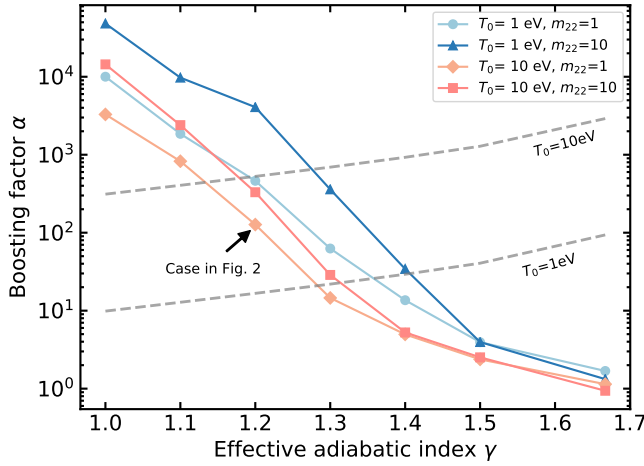


FIG. 3: Strong dependence of the boosting factor α on the effective adiabatic index γ for the four corner cases of (T_0, m_{22}) . α increases with decreasing γ and exceeds two (one) orders of magnitude when $\gamma \lesssim 1.2$ (1.4) regardless of T_0 and m_{22} . For comparison, dashed lines show the Eddington limit with $T_0 = 1$ and 10 eV. The arrow indicates the case in Fig. 2.

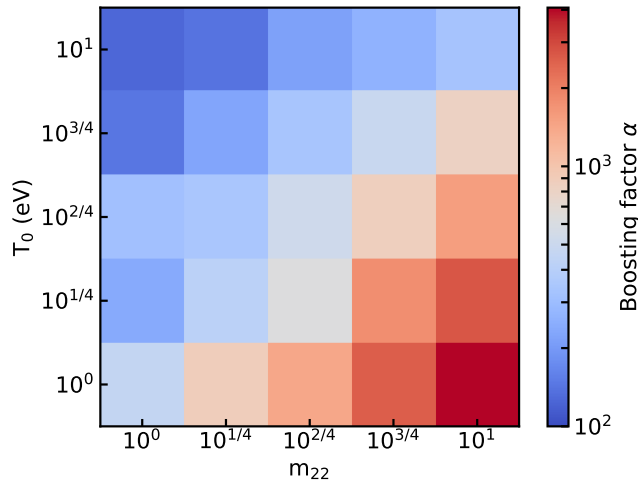


FIG. 4: Boosting factor α as a function of T_0 and m_{22} with $\gamma = 1.2$. α generally exceeds two orders of magnitude and increases with increasing m_{22} and decreasing T_0 .

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