

Dark-matter-independent orbital decay bounds on ultralight bosons from OJ287

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(Received 24 May 2025; accepted 27 October 2025; published 26 November 2025)

Ultralight bosons, predicted in scenarios beyond the Standard Model and viable dark matter (DM) candidates, can form superradiant clouds around spinning black holes influencing their dynamics. Using a century-long monitored OJ287 supermassive black hole binary we set first DM-independent, dynamical constraints on their masses $\mu = (8.5\text{--}22) \times 10^{-22}$ eV. These dynamical constraints, driven by boson cloud friction, are robust against DM-model uncertainties and offer a novel ultralight boson probe. We show that analogous superradiant dynamics across the cosmic population of supermassive black hole systems could help resolve final-parsec evolution stalling problem and imprint a detectable suppression and break in the gravitational wave background.

DOI: [10.1103/y5qn-rvv1](https://doi.org/10.1103/y5qn-rvv1)

I. INTRODUCTION

Ultralight bosons (ULBs) with masses below the eV scale naturally arises in many extensions of the Standard Model (SM) and constitute viable dark matter (DM) candidates (see, e.g., [1–11]). With feeble ULB nongravitational couplings, cosmological and astrophysical probes (see e.g., [12]) can be highly sensitive beyond and complementary to laboratory searches. When the wavelength of ULB is much larger than the black hole (BH) horizon, a rapidly rotating BH can efficiently amplify the bosonic field waves extracting rotational energy and forming long-lived macroscopic “gravitational atom” (GA)

clouds [4,13–15]. This superradiant amplification is the gravitational analog of observed Dicke superradiance in quantum optics and condensed matter [16]. The resulting rich phenomena [17–25] offer intriguing windows for exploring fundamental physics.

In this work we introduce a novel dynamical probe of ULBs based on their superradiance impact on BH binary orbits, independent of any DM assumptions. Unlike previous ULB superradiance constraints derived from BH spin measurement statistics (e.g., [21,26]), our method uses orbital decays as a new dynamical channel to probe the $\mu \sim 10^{-21}$ eV window. We further show that superradiant boson clouds can accelerate supermassive BH binary evolution in the late stage final-parsec regime, where evolution can problematically stall. An ideal system for this test is the OJ287 supermassive BH (SMBH) binary [27], monitored for over a century. Prior analyses of OJ287 crucially relied on new fields constituting DM abundance [27–29]. Here, we set the first DM-independent dynamical limits on ULBs from SMBH orbital evolution, and identify distinctive gravitational wave background signatures, connecting particle physics, astrophysics, and gravitational wave (GW) cosmology.

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II. ULTRALIGHT BOSON SUPERRADIANCE

ULB fields can efficiently seed macroscopic superradiant clouds around spinning BHs when their Compton wavelength $\lambda_c = 1/\mu$, with μ being their mass, exceeds Schwarzschild radius $r_s = 2M_{\text{BH}}$ of BH with mass M_{BH} . As we show, superradiance clouds can significantly impact SMBH binary merger dynamics. We adopt $c = \hbar = G = 1$.

Let us consider a minimal scenario of massive scalar field Φ propagating on a Kerr spacetime background. With Lagrangian contributions $\mathcal{L} \supset -(\partial\Phi)^2/2 - \mu^2\Phi^2/2$, neglecting possible self-interactions (see [13,30–34] for their potential effects). The resulting Klein-Gordon equation governing the ULB behavior can be simplified to a Schrödinger-like equation by expanding it with respect to a small “gravitational fine-structure constant” $\alpha = M_{\text{BH}}\mu < 1$. See Appendix A for details. In this regime, the Schwarzschild radius of the BH is smaller than the Compton wavelength of the ULB and the system can be described as “gravitational atom (GA)” resembling eigenfunctions of a hydrogen atom described by quantum numbers n, l, m as

$$\psi_{nlm}(r, \theta, \phi) \simeq R_{nl}(r)Y_{lm}(\theta, \phi)e^{-i(\omega_{nlm}-\mu)t}, \quad (1)$$

where $R_{nl}(r)$ is the radial wave function and $Y_{lm}(\theta, \phi)$ are the spherical harmonics. The corresponding eigenvalue consists of a real part and an imaginary part $\omega_{nlm} = E_{nlm} + i\Gamma_{nlm}$, which denotes the energy of the eigenstate and its growth or decay rate, respectively. Exponential superradiance growth occurs when $\Gamma_{nlm} > 0$.

The superradiant growth rate of an $|nlm\rangle$ state can be computed analytically using the Detweiler approximation [35], considering $\alpha \lesssim 0.3$ [36,37],

$$\Gamma_{nlm} = 2\tilde{r}_+ C_{nl} g_{lm}(m\Omega_{\text{BH}} - \omega_{nlm})\alpha^{4l+5}, \quad (2)$$

where $\Omega_{\text{BH}} = \chi/2M_{\text{BH}}\tilde{r}_+$ is the angular velocity of the outer BH horizon at $\tilde{r}_+ = 1 + \sqrt{1 - \chi^2}$ and the dimensionless BH Kerr spin parameter is defined as $\chi = J/M_{\text{BH}}^2$, where J being the angular momentum. Here, C_{nl} are numerical coefficients and $g_{lm} = g_{lm}(\chi, \alpha, \mu)$ are functions [15,35]. In a superradiant state with $\Gamma_{nlm} > 0$, ULB cloud will exponentially grow until spin χ becomes insufficient and saturates at

$$\chi_{\text{sat}} = \frac{4m\alpha}{m^2 + 4\alpha^2}. \quad (3)$$

By analogy with electric case, GA Bohr radius is $r_0 = 1/\mu\alpha$, and the excited states have radii $r_n \simeq n^2 r_0$. The density of a boson cloud $\rho_{nlm} = M_{nlm}|\psi_{nlm}|^2$ can be described as (e.g., [38])

$$\rho_{nlm}(x, \theta) = 2.72 \times 10^{12} M_{\odot} \text{pc}^{-3} A_{nlm}(x, \theta) \frac{\beta}{0.1} \left(\frac{\alpha}{0.1}\right)^6 \times \left(\frac{1.83 \times 10^{10} M_{\odot}}{M_{\text{BH}}}\right)^2, \quad (4)$$

where $A_{nlm}(x, \theta)$ denotes the spatial density distribution, where $x = r/r_0$ and β is the mass ratio between the boson cloud and BH. We mainly focus on $|211\rangle$ and $|31-1\rangle$ states in this work,

$$A_{211}(x, \theta) = \frac{1}{64\pi} x^2 e^{-x} \sin^2 \theta, \quad (5)$$

$$A_{31-1}(x, \theta) = \frac{1}{6561\pi} (x^2 - 6x)^2 e^{-2x/3} \sin^2 \theta. \quad (6)$$

The $|211\rangle$ mode is the fastest growing, as we discuss in Appendix A. Moreover, as shown in Appendix B, the $|211\rangle$ mode can transform to the $|31-1\rangle$ mode that is then slowly absorbed by the BH. In the following discussion, we consider the impact of superradiant bosons in OJ287 orbital dynamics as illustration in Fig. 1. The density profile of superradiant bosons plays an important role in determining the orbital dynamics of OJ287, we display the corresponding profiles and enclosed cloud mass for parameters relevant to OJ287 ($\alpha = \beta = 0.1$) in Fig. 2.

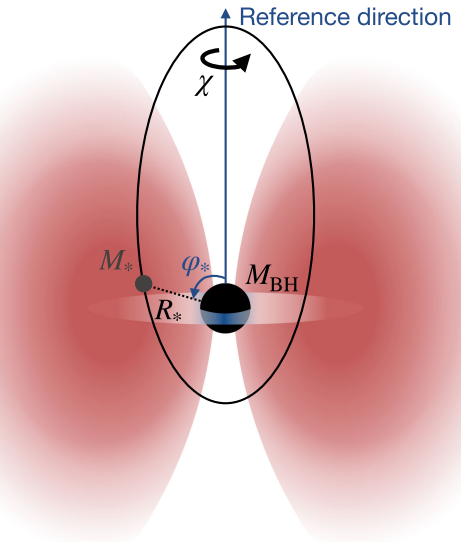


FIG. 1. Schematic illustration of secondary BH (gray) of mass M_* (black) traversing ULB superradiance cloud in $|211\rangle$ state orbiting primary BH of mass M_{BH} and spin χ in a SMBH binary, such as the OJ287 system. Primary BH rotation defines the reference axis. The secondary's position is given by its separation R_* and orbital phase ϕ_* measured from this axis. The orbital angular momentum is perpendicular to the primary BH spin.

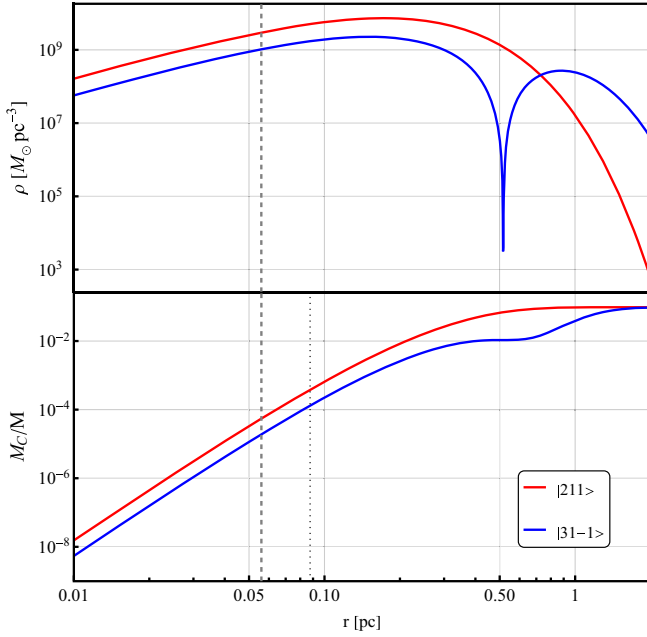


FIG. 2. [Top] Density profiles of the $|211\rangle$ and $|31-1\rangle$ GA states for $\alpha = \beta = 0.1$ and $\theta = \pi/2$ around primary BH in the OJ287 system. The corresponding Bohr radius is $r_0 = 0.088$ pc. [Bottom] Ratio of enclosed superradiance cloud mass relative to SMBH M_{BH} . The vertical dashed line mark the semimajor axis (0.056 pc) of OJ287. The dotted line corresponds to $50r_s$.

III. OJ287 DYNAMICS AND NOVEL BOSON LIMITS

A pristine SMBH system to explore the dynamical effects of superradiance clouds is blazar OJ287. OJ287 is an unusually massive and tight BH binary with primary BH mass of $M_{\text{BH}} = 1.83 \times 10^{10} M_{\odot}$, secondary BH mass of $M_* = 1.50 \times 10^8 M_{\odot}$ and orbital timing measurements collected over ~ 120 yr period. Its orbital parameters include semi-major axis $a = 0.056$ pc, inclination angle $i = 90^\circ$, and eccentricity $e = 0.657$ [39–41] as listed in Table I. The binary is shrinking due to continuous GW emission expected from general relativity [42],

TABLE I. Parameters characterizing OJ287 BH system.

Parameters	Value	Ref.
Primary SMBH mass ($10^{10} M_{\odot}$)	1.8348 ± 0.0008	[39]
Secondary SMBH mass ($10^8 M_{\odot}$)	1.5013 ± 0.0025	[39]
Eccentricity e	0.657 ± 0.001	[39]
Precession angle per period $\Delta\vartheta$	$38.62^\circ \pm 0.01^\circ$	[39]
Orbital period T (yr)	12.067 ± 0.007	[39]
Orbital period decay rate $\dot{T}$	-0.00099 ± 0.00006	[39]
Spin of the primary SMBH χ_{P_2}	0.381 ± 0.004	[39]
Spin of the primary SMBH χ_{P_1}	0.313 ± 0.01	[43]
Semi-major axis a (pc)	0.056	[40]
Inclination angle of orbit plane	90°	[41]

$$\langle P_{\text{GW}} \rangle = -\frac{32m_r^2 M_t^3}{5a^5} (1-e^2)^{-7/2} \left(1 + \frac{73}{24}e^2 + \frac{37}{96}e^4 \right), \quad (7)$$

where $m_r = M_{\text{BH}} M_* / (M_* + M_{\text{BH}})$ is the reduced system mass and $M_t = M_* + M_{\text{BH}}$ is the total mass.

Recent work claimed an apparent OJ287 total orbital-decay power $\langle P_{\text{tot}} \rangle = (3.66 \pm 0.24) \times 10^{41}$ W seemingly exceeds the quadrupole GW prediction $\langle P_{\text{GW}} \rangle = (2.62 \pm 0.02) \times 10^{41}$ W by $\sim 4.3\sigma$ [29]. However, detailed reanalysis finds only a 0.41σ deviation [27]. Remaining conservative, we base our limits on the null-excess result compared to relativity GW predictions and derive the first DM-independent, dynamical constraint on ULBs from OJ287.

ULBs of mass $\mathcal{O}(10^{-21})$ eV can form a superradiance cloud around a primary SMBH of OJ287 within the current age of the Universe $\tau_U \simeq 13.8$ Gyr, considering high initial spin χ_i as expected of most SMBHs [44]. This can be seen from the superradiance growth rate of the fastest growing GA state $|211\rangle$,

$$\Gamma_{211} \simeq 1.2 \times 10^{-16} \left(\frac{\alpha}{0.1} \right)^9 \left(\frac{1.8 \times 10^{10} M_{\odot}}{M_{\text{BH}}} \right) \text{s}^{-1}. \quad (8)$$

The $|211\rangle$ cloud has a typical Bohr radius of $r_0 \simeq \mathcal{O}(0.01-1)$ pc and density of ρ_{211} . Hence, the $|211\rangle$ GA state can be fully saturated and with the cloud mass of $M_{211} \simeq \alpha M_{\text{BH}}$ [45]. The total cloud growth time takes number of e -folds that can be up to $\log(\Delta\chi M_{\text{BH}}^2) \sim 170$ times larger than $1/\Gamma_{211}$, which is still significantly less than τ_U for OJ287.

The superradiant cloud could potentially deplete through variety of channels, such as monochromatic GW emission [46,47] or Landau-Zener bound-to-bound resonances [22,23,48–50] and bound-to-unbound resonance, known as ionization [24,51]. Detailed analysis demonstrates (see Appendix B for details) that the OJ287 system remains not significantly affected even when these effects are taken into account.

As the secondary BH of mass M_* traverses the superradiance GA cloud of local density ρ , it experiences a gravitational drag force acting as dynamical friction [52], which in the nonrelativistic, wavelike regime reads [53]

$$F_{\text{DF}} = \frac{4\pi M_*^2 \rho}{v^2} C_{\Lambda}(\xi, kr_{\Lambda}), \quad (9)$$

where v is the speed relative to the boson cloud that equals $\alpha\sqrt{(1+q)/x_*}$ in circular orbit and $\alpha\sqrt{(1+q)(2/x_* - r_0/a)}$ in elliptic orbit with $x_* = R_*/r_0$, $\xi = M_*\mu/v$, $k = \mu v$, and r_{Λ} representing the smaller quantity between the size of the orbit and the size of the cloud. In the $\xi \ll 1$ limit the Coulomb logarithm becomes

$$C_\Lambda(kr) = \text{Cin}(2kr) + \frac{\sin 2kr}{2kr} - 1, \quad (10)$$

with

$$\text{Cin}(z) = \int_0^z (1 - \cos t) \frac{dt}{t}. \quad (11)$$

Writing the instantaneous frictional power as $P_{\text{DF}} = F_{\text{DF}} v$, and $q = M_*/M_{\text{BH}}$ being the binary mass ratio, one finds

$$P_{\text{DF}} = 4\pi \frac{q^2 M_{\text{BH}}^2 \rho(x_*, \theta_*)}{\alpha \sqrt{(1+q)(2/x_* - r_0/a)}} C_\Lambda(kr_\Lambda). \quad (12)$$

For eccentric orbits with semimajor axis a , eccentricity e , and eccentric anomaly z with $r = a[1 - e \cos z]$, the effective impact parameter kr_Λ is given by [54]

$$kr_\Lambda = \begin{cases} \sqrt{x_p} \frac{(1-e \cos z)^{3/2}}{\sqrt{1+e \cos z}}, & x_p \leq x_{97}, \\ \frac{x_{97}}{\sqrt{x_p}} \frac{(1-e \cos z)^{3/2}}{\sqrt{1+e \cos z}}, & x_p > x_{97}, \end{cases} \quad (13)$$

where $x_p = r_p/r_0$ is the periapsis in units of r_0 , and x_{97} encloses 97% of the cloud mass.

To estimate the impact of boson cloud-induced dynamical friction on SMBH binary evolution, we compute dynamical friction power averaged over one period. The orbit of SMBH binary can be described as

$$R_*(\varphi_*) = \frac{a(1-e^2)}{1+e \cos(1-\zeta)\varphi_*}, \quad (14)$$

where ζ is the precession phase angle per period. When a secondary SMBH companion is traversing the GA boson cloud, it experiences position-dependent cloud density $\rho(R_*, \theta_*)$. The angular position θ_* depends on the orbital phase φ_* and inclination angle ι as $\cos \theta_* = \cos \varphi_* \sin \iota$. The detailed OJ287 parameters are listed in Table I.

Then, the average of Eq. (12) over one orbital period T is

$$\langle P_{\text{DF}} \rangle = \frac{1}{T} \int_0^{2\pi} P_{\text{DF}}(R_*(\varphi_*), \theta_*(\varphi_*, \iota)) \frac{dt}{d\varphi_*} d\varphi_*. \quad (15)$$

Here,

$$\frac{d\varphi_*}{dt} = \frac{\sqrt{(1+q)M_{\text{BH}}a(1-e^2)}}{R_*(\varphi_*)^2} \quad (16)$$

and the orbital period T can be calculated as

$$T = \int_0^{2\pi} \frac{dt}{d\varphi_*} d\varphi_*. \quad (17)$$

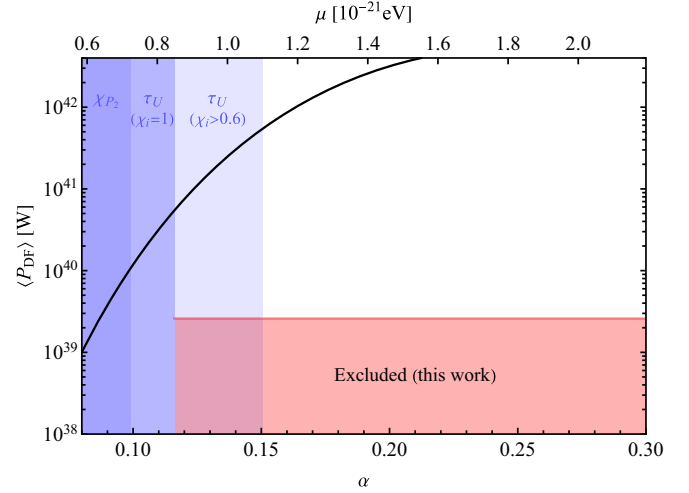


FIG. 3. The black solid line shows the orbital decay power from dynamical friction of dominant superradiance cloud state in OJ287. The red region denotes the upper limit on the excess orbital decay power, consistent with general relativity GW predictions $\langle P_{\text{DF}} \rangle / \langle P_{\text{GW}} \rangle < 0.01$ [27]. Boson masses μ for which the black curve rises above the red region produce excessive decay and are therefore excluded, shown by horizontal span of the red region. The region where saturated spin is below the measured primary BH spin value χ_{P_2} (dark blue) is shown. Light blue regions show the parameters with which the boson cloud cannot form within cosmic time τ_U , considering different initial spins χ_i .

We employ the resulting Eq. (15) to compute the orbital decay rate due to ULB clouds. In Fig. 3 we display computed orbital-decay power generated by dynamical friction in OJ287 due to a saturated GA $|211\rangle$ state. While our baseline analysis assumes no excess beyond emission predicted by relativity, our framework can also accommodate a genuine deviation if such is confirmed. For GA couplings $\alpha = 0.121\text{--}0.127$, corresponding to boson masses $\mu = (8.8\text{--}9.3) \times 10^{-22}$ eV, the predicted dissipation can accommodate claimed excess in orbital decay power $\langle P_{\text{DF}} \rangle = \langle P_{\text{tot}} \rangle - \langle P_{\text{GW}} \rangle$ of [29]. In our analysis we conservatively adopt null-excess [27], and derive new DM-independent bounds. Here, the superradiance growth time obeys $170 \times \Gamma_{211}^{-1} < \tau_U$ and the primary spin satisfies $\chi < \chi_{211,\text{sat}}$, with χ_{P_2} being more conservative compared to considering χ_{P_1} . Hence, the $|211\rangle$ state is ensured to be populated and stable. See Appendix B for further details.

Throughout this band the superradiance growth time obeys $170 \times \Gamma_{211}^{-1} < \tau_U$ and the primary spin satisfies $\chi < \chi_{211,\text{sat}}$, with χ_{P_2} being more conservative compared to considering χ_{P_1} . Hence, the $|211\rangle$ state is ensured to be populated and stable. See Appendix B for further details.

Our scenario is consistent with timing observations of OJ287 flares, which are thought to originate from secondary BH interacting with the primary BH accretion disk. Integrating the density profile for the $|211\rangle$ state out to $50r_s$

yields a cloud mass of $\sim 0.05\%$ of the primary BH mass as shown in bottom panel of Fig. 2. This is well below the $\sim 3\%$ mass ratio at which the flare timing predictions could be impacted, suggested from [28].

We set novel limits on ULBs independent of assumptions on their DM modeling and cosmological abundance considering conservatively no excess of orbital-decay power over relativity GW predictions at 0.41σ at 1PN level [27], which corresponds with observational limit power ratio $\langle P_{\text{DF}} \rangle / \langle P_{\text{GW}} \rangle \lesssim 0.01$ in Fig. 3. Hence, this forbids ULB masses larger than 6.4×10^{-22} eV considering superradiance effects. Combining with requirements that superradiant boson cloud efficiently forms satisfying $170 \times \Gamma_{211}^{-1} < \tau_{\text{U}}, \chi < \chi_{211, \text{sat}}$, and Detweiler approximation validity range $\alpha \lesssim 0.3$, we set novel bounds restricting the ULB mass range $\mu \simeq (8.5\text{--}22) \times 10^{-22}$ eV. Since additional model-dependent effects such as ULB self-interactions only suppress superradiant growth, our derived limits are conservative. Varying OJ287 parameters within 1σ uncertainties that we consider excluded ULB limit region is not significantly affected.

Our ULB constraints are the first derived from OJ287 orbital dynamics that are independent of DM-assumptions, unlike previous studies [27–29]. These studies either attempt to explain the orbital evolution of the OJ287 system through the presence of a DM overdensity spike or use the OJ287 observations to constrain such a spike, which requires a substantial concentration of DM around the central black hole for its formation. In contrast, the formation of a superradiant boson cloud we consider does not require the ULB to constitute a significant fraction of the DM density, since the superradiance instability allows even a small bosonic population to efficiently extract rotational energy from the black hole and form a macroscopic condensate. Therefore, our result provides a DM-independent constraint on the ULB mass.

IV. FINAL PARSEC EVOLUTION

GA superradiance clouds stemming from ULBs of mass around $\mu \simeq \mathcal{O}(10^{-21})$ eV, as suggested by OJ287 orbital decay analysis, envelop $\mathcal{O}(10^{10})M_{\odot}$ SMBHs with Bohr radii $r_c \sim \mathcal{O}(1)$ pc. At this regime the evolution of SMBH binaries could stall, known as the “final-parsec problem”, denoting inefficient SMBH binary passage through its final parsec separation within a Hubble timescale [55]. Recently reported pulsar-timing arrays (PTAs) detection of a stochastic GW background in the nanohertz band [56–59] that can originate from SMBHs further highlights these effects [60]. Hence, the dynamical friction from GA boson clouds in SMBH binaries could efficiently accelerate the shrinkage of binary separation within final parsec evolution.

We estimate that bound-to-bound resonance transitions deplete at most $\mathcal{O}(10)\%$ of the $|211\rangle$ population for equal-mass binaries and $\lesssim 10\%$ for $q = 0.1$. See Fig. 6 for further

details. This is distinct from OJ287 where the GA is taken to be almost purely in the $|211\rangle$ state. To capture this behavior we model the GA cloud distribution as $\rho = \eta\rho_{211} + (1 - \eta)\rho_{31-1}$, with η denoting relative contribution of $|211\rangle$ and $|31-1\rangle$ states that depends on SMBH binary parameters. For $q = 0.1$ binaries we have $\eta \simeq 0.9$, while for $q = 1$ case $\eta \simeq 0.1$. Here, we estimate negligible contributions of other states, although our approach can be readily extended to incorporate this. Then the orbital decay rate of SMBH binaries can be found by semimajor axis evolution as

$$\frac{1}{a} \frac{da}{dt} = \frac{2a}{qM_{\text{BH}}^2} (\langle P_{\text{DF}} \rangle + \langle P_{\text{GW}} \rangle). \quad (18)$$

The resulting elapsed time within final parsec τ_{pc} can be estimated by integrating Eq. (18) over evolving separations to obtain its evolution time as follows:

$$\tau = \int \frac{qM_{\text{BH}}^2}{\langle P_{\text{DF}} \rangle + \langle P_{\text{GW}} \rangle} \frac{1}{2a^2} da. \quad (19)$$

Solving the final-parsec problem requires the evolution time τ smaller than 1 Gyr, and this put constraints on SMBH mass and boson mass. We consider the evolving separation from 2 pc to 0.01 pc in the top panel of Fig. 4. Here, we partially release this evolving separation from 0.5 pc to 0.1 pc, and the resulting allowed parameter regions are shown in the bottom panel of Fig. 4. It shows the existence of superradiant boson cloud can help $\mathcal{O}(10^8\text{--}10^{10})M_{\odot}$ SMBH binaries overcome their final $\mathcal{O}(0.1)$ pc separations.

In Fig. 4 we display the region in the (M_{BH}, μ) plane for which the inspiral during final-parsec evolution is completed in $\tau_{\text{pc}} < 1$ Gyr. Here, the lower curve denotes $\Gamma_{211}^{-1} < \tau_{\text{U}}$, and with dashed line showing conservative $170 \times \Gamma_{211}^{-1} < \tau_{\text{U}}$. Since the parameters relevant for OJ287 ($1.8 \times 10^{10}M_{\odot}, 0.9 \times 10^{-21}$ eV) lie inside the allowed band this indicates that dynamical friction from the ULB cloud can drive the binary through the final parsec before GW backreaction takes over, and can simultaneously account for the claimed OJ287 orbital decay excess if confirmed. Here, we also display for reference the circular binary case with eccentricity $e = 0$. Systems with ULB clouds that satisfy the superradiance constraint but reside outside of the shaded region will also partially accelerate the decay. This indicates that our scenario can alleviate the final-parsec problem in various SMBH systems.

V. GRAVITATIONAL WAVE BACKGROUND AND TURNOVER

Once a SMBH binary evolution has crossed the final parsec, GW emission dominates and sources nanohertz stochastic GW background observable by PTAs.

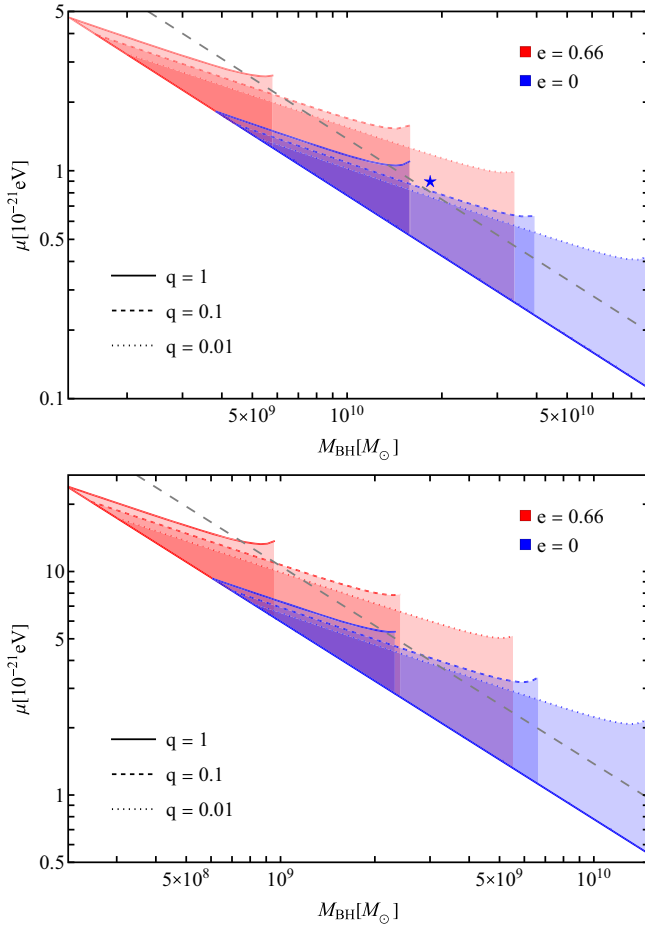


FIG. 4. [Top] Allowed SMBH mass and ULB mass parameter space where binary evolution time from 2 pc to 0.01 pc is shorter than $\tau_{\text{pc}} < 1$ Gyr, alleviating the final-parsec problem. The gray dashed line corresponds to conservative superradiance timescales of $170 \times \Gamma_{211}^{-1} < \tau_U$. The parameters for addressing OJ287 orbital decay ($1.83 \times 10^{10} M_\odot$, 0.9×10^{-21} eV) are shown by a blue star. [Bottom] Allowed SMBH mass and ULB mass parameter space where binary evolution time from 0.5 pc to 0.1 pc is shorter than $\tau_{\text{pc}} < 1$ Gyr, alleviating the final-parsec problem.

A superradiant cloud drag reduces the energy radiated per logarithmic frequency interval,

$$\frac{dE_{\text{GW}}}{df_s} = \frac{\pi^{2/3} M_{\text{BH}}^{5/3}}{3 f_s^{1/3}} \frac{q}{(1+q)^{1/3}} \frac{\langle P_{\text{GW}} \rangle}{\langle P_{\text{GW}} \rangle + \langle P_{\text{DF}} \rangle}, \quad (20)$$

with a characteristic strain

$$h_c^2(f) = \frac{4}{\pi f} \int dz dq dM_{\text{BH}} \frac{d^3 n}{dz dq dM_{\text{BH}}} \frac{dE_{\text{GW}}}{df_s}, \quad (21)$$

where $f_s = f(1+z)$ is the source frequency and $d^3 n/dz dq dM_{\text{BH}}$ is the SMBH binary population that we estimate considering galaxy merger rate distribution, where we follow [61,62] to evaluate SMBH population (see

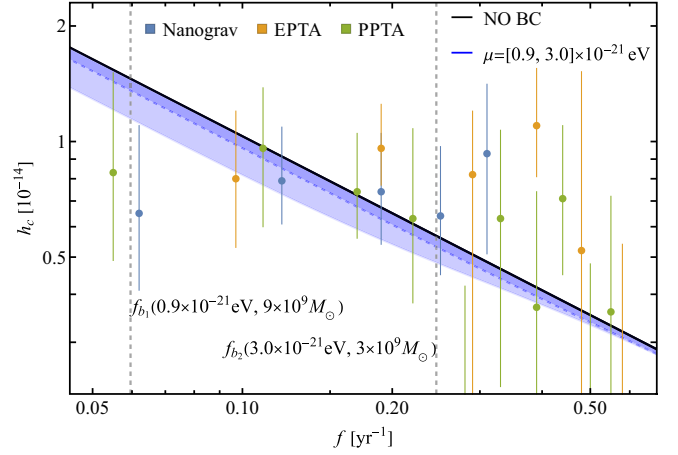


FIG. 5. GW background strain resulting from cosmic population of SMBH binaries considering ULB cloud dynamical friction effects for a range of boson masses. Black lines denotes standard relativity GW predictions. Shaded region above the blue solid line depicts superradiance effects with all SMBHs having maximal initial spin, and shaded region above blue dashed line depicts conservative case of $170 \times \Gamma_{211}^{-1} < \tau_U$. The vertical lines show new turnover GW frequency signatures considering f_b for different (μ, M_{BH}) cases. Observations from NANOGrav [56], EPTA [64], and PPTA [65] are shown.

Appendix C). In evaluating Eq. (20) we account for distinct GA configurations by appropriately considering cloud distribution profile ρ with ratio of states η set depending on SMBH binary and ULB parameters. With superradiant growth relevant when $\Gamma_{211}^{-1} < \tau_U$, we approximate that only SMBHs above a critical mass M_{crit} host GA clouds for a given μ . Hence, dynamical friction acts for $M_{\text{BH}} \gtrsim M_{\text{crit}}$ in the mass integral range 10^5 – $10^{10} M_\odot$. We denote with band parameter range extending to conservative $170 \times \Gamma_{211}^{-1} < \tau_U$. This includes contributions from partially populated boson clouds. Since significant fraction of SMBHs is expect to have large spin [44], we consider a scenario where all SMBHs are initial maximal spin, although our results can be readily linearly rescaled for smaller fraction.

In Fig. 5 we show that the GA cloud effects of our scenario can suppress the PTA strain by $\mathcal{O}(10\text{--}30)\%$ relative to the vacuum prediction, improving agreement with the observed PTA amplitude, where the predicted strain in GA cloud scenario gives $\chi^2 = 16.5$ (blue solid—all SMBHs with maximal spin), $\chi^2 = 17.5$ (blue dashed—conservative superradiance timescales), while predicted strain in GW scenario gives $\chi^2 = 18.8$ (black solid—general relativity prediction). (see App. D for statistical analysis). For completeness we note that saturated GA clouds continuously radiate in GWs [63]. This lies in ~ 0.1 microhertz band for 10^{-21} eV ULBs and contributes negligibly to the nanohertz background.

Since the superradiant cloud drag becomes ineffective once the SMBH orbit shrinks inside its Bohr radius r_0 , GW

background exhibits a turnover at

$$f_b \simeq 3.2 \text{ nHz} \left(\frac{\mu}{10^{-21} \text{ eV}} \right)^3 \left(\frac{M_{\text{BH}}}{10^{10} M_\odot} \right)^2. \quad (22)$$

This is a distinct testable feature of our scenario, distinguishing it from other proposals such as those based on DM spikes (e.g., [29,62]) and solitonic cores [66] (see Appendix E) that rely on assumptions about DM abundance. Detecting or ruling out the predicted turnover with upcoming precision PTA observations would thus isolate the dynamical imprint of superradiant clouds and pin down ULB mass scales.

Our results connect particle physics, astrophysics, GW astronomy and cosmology. In particle physics, the derived ULB mass bounds inform future searches beyond the DM paradigm. In astrophysics, boson cloud drag accelerates SMBH coalescence linking to the final-parsec evolution problem. In GW astronomy and cosmology, the predicted GW background suppression and turnover is directly testable with upcoming PTA data.

VI. CONCLUSIONS

We have introduced a novel DM-independent dynamical probe of ULBs based on orbital dynamics. Applying it to a century-monitored OJ287 system, no excess over relativity predictions excludes ULBs in the mass range $(8.5\text{--}22) \times 10^{-22} \text{ eV}$ setting the first DM-independent ULB dynamical constraint derived from binary BH dynamics rather than spin statistics. Our framework enables accelerating supermassive-BH binaries through the long standing final-parsec evolution bottleneck. Consequent suppression and turnover of the nanohertz gravitational wave background provide a decisive test for forthcoming PTA and high-precision orbital monitors. Our framework thus links distinct fields of particle physics, BH astrophysics, and gravitational wave cosmology in a single observational channel.

ACKNOWLEDGMENTS

We thank Masha Baryakhtar and Yong Tang for comments. H. Y. Z. would like to thank Sida Lu and Yi Wang for the fruitful discussion on final-parsec problem. Q. D. and H. Y. Z. were supported by IBS under the Project Code No. IBS-R018-D3. M. H. was supported by IBS under the Project Code No. IBS-R018-D1. V. T. acknowledges support by the World Premier International Research Center

Initiative (WPI), MEXT, Japan and JSPS KAKENHI Grant No. 23K13109.

DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: SCALAR-FIELD SUPERRADIANT GRAVITATIONAL ATOM

A scalar field of mass μ in the Kerr background of a rotating BH of mass M_{BH} and dimensionless spin χ obeys the Klein-Gordon equation,

$$(\square_{\text{Kerr}} - \mu^2)\Phi = 0, \quad (\text{A1})$$

where $\square_{\text{Kerr}}$ is the d'Alembertian operator associated with the Kerr metric. With appropriate boundary conditions, Eq. (A1) separates into Teukolsky radial and angular equations describing perturbations of a Kerr BH. This can be solved numerically with Leaver's method [67,68]. We consider analytical solutions in the nonrelativistic limit.

For the gravitational fine-structure constant $\alpha = \mu M_{\text{BH}} \ll 1$ one can factor out the fast oscillating phase,

$$\Phi(t, \mathbf{x}) = \frac{e^{-i\mu t}}{\sqrt{2\mu}} \psi(t, \mathbf{x}) + \text{c.c.}, \quad (\text{A2})$$

such that Eq. (A1) reduces at leading order in α to the hydrogen atom analog Schrödinger-like equation describing the “gravitational atom (GA)”,

$$i\partial_t \psi = \left(-\frac{\nabla^2}{2\mu} - \frac{\alpha}{r} + \mathcal{O}(\alpha^2) \right) \psi. \quad (\text{A3})$$

Then, the bound GA states are labeled by quantum numbers (n, ℓ, m) with wave functions

$$\psi_{n\ell m} = R_{n\ell}(r) Y_{\ell m}(\theta, \phi) e^{-i(\omega_{n\ell m} - \mu)t}, \quad (\text{A4})$$

where $R_{n\ell}$ denotes radial functions and $Y_{\ell m}$ spherical harmonics.

The GA eigenfunctions are described by complex eigenfrequencies $\omega_{n\ell m} = E_{n\ell m} + i\Gamma_{n\ell m}$. With the Detweiler approximation [35], considering $\alpha \lesssim 0.3$ [36,37], the real and imaginary components are analytically expressed as

$$\begin{cases} E_{n\ell m} = \mu \left[1 - \frac{\alpha^2}{2n^2} - \frac{\alpha^4}{8n^4} - \frac{(3n-\ell-1)\alpha^4}{n^4(\ell+1/2)} + \frac{2\chi m \alpha^5}{n^3 \ell(\ell+1/2)(\ell+1)} + \mathcal{O}(\alpha^6) \right], \\ \Gamma_{n\ell m} = 2\tilde{r}_+ C_{n\ell} g_{\ell m} (m\Omega_{\text{BH}} - \omega_{n\ell m}) \alpha^{\ell+5}, \end{cases} \quad (\text{A5})$$

where $\tilde{r}_+$ denotes outer horizon of the BH, C_{nl} are numerical coefficients, and $g_{lm} = g_{lm}(\chi, \alpha, \mu)$ are functions defined in e.g., [15,35,69]. Exponential superradiant growth occurs when condition $m\Omega_{\text{BH}} > \omega_{n\ell m}$ is satisfied, which happens first for the fastest-growing mode $|211\rangle$. Higher GA ℓ -modes grow significantly more slowly with $\Gamma \propto \alpha^{4\ell+5}$ since $\alpha \ll 1$. Hence, the BH retains finite spin over cosmic times.

The cloud itself can slowly deplete via GW emission at a rate,

$$\Gamma_{\text{GW}}(n, l, m) = -B_{nl} \frac{S_C/m}{M^2} \mu \alpha^{4l+10}, \quad (\text{A6})$$

where S_C is the cloud angular momentum and numerical coefficient B_{nl} can be found in [17]. For the fastest growing $|211\rangle$ the emission rate is around $\Gamma_{\text{GW}} \simeq -2.11 \times 10^{-22} \text{s}^{-1}$. Hence, the respective depletion time exceeds the age of the universe τ_U .

The cloud density $\rho_{nlm} = M_{nlm} |\psi_{nlm}|^2$ can be written as

$$\rho_{nlm}(r, \theta) = \beta A_{nlm}(x, \theta) \frac{M_{\text{BH}}}{r_0^3}, \quad (\text{A7})$$

with $x = r/r_0$, $r_0 = (\mu\alpha)^{-1}$ being the Bohr radius, and $\beta = M_{nlm}/M_{\text{BH}}$.

APPENDIX B: IMPACT OF BINARY COMPANION ON GRAVITATIONAL ATOMS

In BH binaries, boson cloud states could be affected by a companion BH.

A companion of mass M_* at an orbital separation $R_*(t)$ can perturb the boson cloud gravitationally. In the reference frame of central BH, considering $\vec{r} = \{r, \theta, \phi\}$ and keeping only the Newtonian contributions of the companion's potential, the gravitational perturbation due to binary companion is

$$V_*(t, \vec{r}) = -\alpha q \sum_{\ell_* \geq 2} \sum_{|m_*| \leq \ell_*} \mathcal{E}_{\ell_* m_*}(\iota_*, \varphi_*) Y_{\ell_* m_*}(\theta, \phi) F_{\ell_*}(r), \quad (\text{B1})$$

where

$$F_{\ell_*}(r) = \frac{r^{\ell_*}}{R_*^{\ell_*+1}} \Theta(R_* - r) + \frac{R_*^{\ell_*}}{r^{\ell_*+1}} \Theta(r - R_*) \quad (\text{B2})$$

and Θ is the Heaviside step-function, $q = M_*/M_{\text{BH}}$ is the mass ratio of the companion and the GA. Here we have also introduced the tidal moments,

$$\mathcal{E}_{\ell_* m_*} = \frac{4\pi}{2\ell_* + 1} Y_{\ell_* m_*}^*(\theta_*, \phi_*). \quad (\text{B3})$$

The angle θ_* represents the angle between the BH spin direction and the companion location, while ϕ_* represents

the projection of the angular phase of the companion orbit with respect to the primary BH onto the BH equatorial plane. In the Fermi normal reference frame the dipole contribution ($\ell_* = 1$) vanishes [22,69], although [24] claims that a residual dipole might exist. Here, we focus on contributions starting with the quadrupole ($\ell_* = 2$).

The companion perturbation can induce mixing of different GA eigenstates. Considering boson cloud states $|n\ell m\rangle$ and $|n'\ell' m'\rangle$ the matrix element factorizes,

$$\langle n'\ell' m' | V_* | n\ell m \rangle = (-1)^{m'+1} \alpha q \sum_{\ell_* m_*} \mathcal{E}_{\ell_* m_*} \mathcal{G}_{-m' m_* m}^{\ell' \ell_* \ell} I_r, \quad (\text{B4})$$

where I_r represents the radial integral

$$I_r = \int_0^{R^*} r^2 dr R_{n'\ell'}^*(r) R_{nl}(r) \frac{r^{l_*}}{R_*^{l_*+1}} + \int_{R^*}^{\infty} r^2 dr R_{n'\ell'}^*(r) R_{nl}(r) \frac{R_*^{l_*}}{r^{l_*+1}}. \quad (\text{B5})$$

Here, $\mathcal{G}_{-m' m_* m}^{\ell' l_* l}$ is the Gaunt integral

$$\mathcal{G}_{-m' m_* m}^{\ell' l_* l} = \int d\Omega Y_{\ell' -m'}(\Omega) Y_{l_* m_*}(\Omega) Y_{lm}(\Omega). \quad (\text{B6})$$

Requiring that Gaunt integral does not vanish imposes that the following selection rules hold; $-m' + m_* + m = 0$, $|\ell - \ell'| \leq \ell_* \leq \ell + \ell'$, and $\ell + \ell_* + \ell' = 2p$, $p \in \mathbb{Z}$.

With the Newtonian potential, we can construct the new Hamiltonian $H = H_0 + V_*$ consisting of a free part,

$$H_0 = -\frac{1}{2\mu} \partial_{\mathbf{r}}^2 - \frac{\alpha}{r} + \mathcal{O}(\alpha^2), \quad (\text{B7})$$

and a tidal perturbation part V_* . The existence of V_* leads to three phenomena: direct off-resonance coupling of eigenmodes, Landau-Zener (LZ) resonance transitions, and ionization.

For illustration, we consider two cloud states $|1\rangle$ and $|2\rangle$ mixed by the tidal potential V_* of the companion. We define the state coupling as $\eta(t) = V_{21} \langle 2 | V_* | 1 \rangle$ and take $m_{1,2}$ to be the azimuthal quantum numbers of the two levels. The Hamiltonian in the interaction basis is

$$H = \begin{pmatrix} \omega_1 + V_{11} & V_{12} \\ V_{21} & \omega_2 + V_{22} \end{pmatrix} = \begin{pmatrix} \bar{E}_1(t) + i\Gamma_1 & \eta^*(t) \\ \eta(t) & \bar{E}_2(t) + i\Gamma_2 \end{pmatrix}, \quad (\text{B8})$$

where $\bar{E}_{1,2}$ and $\Gamma_{1,2}$ are the unperturbed energy and growth rates, respectively.

Transforming to the frame corotating with the companion gives the leading modification of the superradiant rate of the growing level [70,71]. Taking $|1\rangle$ to represent the superradiant state such as $|211\rangle$ and $|2\rangle$ to

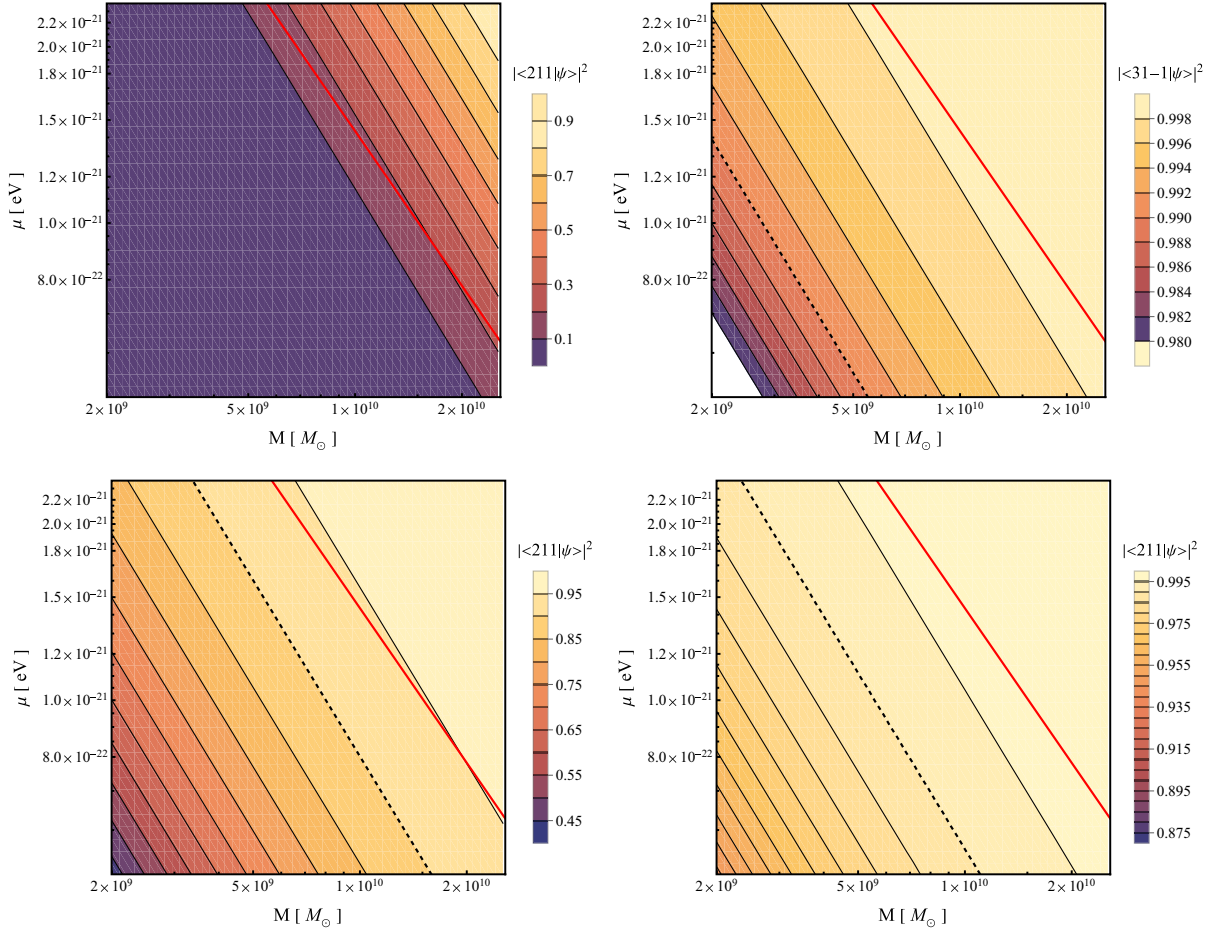


FIG. 6. Survival probability of the GA superradiant eigenmodes under successive resonances. The red curve marks $170 \times \Gamma_{211}^{-1} = \tau_U$, above which the $|211\rangle$ cloud can form within a Hubble time. [Top left] $|211\rangle \rightarrow |31-1\rangle$ transition at $q = 1$. [Top right] $|31-1\rangle \rightarrow |62-2\rangle$ at $q = 1$, with the black dashed line showing the 99% survival contour. [Bottom left] $|211\rangle \rightarrow |31-1\rangle$ at $q = 0.1$. The dashed curve is the 90% contour. [Bottom right] $|211\rangle \rightarrow |31-1\rangle$ with the mass ratio of OJ287 binary, with the 99% survival contour.

represent the absorptive state such as $|21-1\rangle$, the instantaneous correction to superradiant rate of state $|1\rangle$ is

$$\Delta\Gamma_1 \simeq -\frac{\Gamma_1 - \Gamma_2}{[(E_1 - E_2) - (m_1 - m_2)\dot{\varphi}_*(t)]^2} |\eta(t)|, \quad (\text{B9})$$

which is always negative and largest near resonance. For a general binary orbit and eccentricity, one should average over a period [72],

$$\langle \Delta\Gamma_1 \rangle = \frac{1}{T} \int_0^{2\pi} \Delta\Gamma_1 \frac{d\varphi_*}{\dot{\varphi}_*}, \quad (\text{B10})$$

which shows that the effect scales with $|\eta|$.

In the case of OJ287, the orbital angular momentum is nearly orthogonal to the BH spin, i.e., inclination angle $i \simeq 90^\circ$ (see schematic Fig. 1). This strongly suppresses the overlap $|\eta|$ and correction to the superradiant rate. We estimate $\langle \Delta\Gamma_1 \rangle / \Gamma_1 \ll 1$. Thus, off-resonant mixing can be safely neglected in our analysis.

The LZ resonance transitions constitute other prominent effects. We observe that Eq. (B9) develops a pole when the driving frequency $\dot{\varphi}_*$ sweeps through $\Omega_r = (E_1 - E_2)/(m_1 - m_2)$. Close to the crossing we linearize the chirp, $\dot{\varphi}_*(t) \simeq \dot{\varphi}_*(t_0) + \gamma t$ where $\gamma = \gamma_{\text{GW}} + \gamma_{\text{DF}} + \gamma_{\text{BR}}$ takes into account the energy loss via GW, dynamical friction and backreaction of resonance effect on the companion. Here γ_{GW} can be calculated Eq. (2.28) in [22] and γ_{DF} can be estimated by $\langle P_{\text{DF}} \rangle / \langle P_{\text{GW}} \rangle \times \gamma_{\text{GW}}$. Since the backreaction effect is subdominant, we ignore the γ_{BR} in our calculation. and treat the two-level system with the standard LZ formalism [22,73]. With initial populations $c_1(-\infty) = 1$, $c_2(-\infty) = 0$ the occupation after the passage is

$$|c_1|^2 = e^{-2\pi Z}, \quad |c_2|^2 = 1 - e^{-2\pi Z}, \quad Z = \frac{|\eta|^2}{\gamma}, \quad (\text{B11})$$

where Z is the usual LZ parameter.

For a saturated $|211\rangle$ cloud three LZ resonances are in principle relevant; $|211\rangle \rightarrow |31-1\rangle \rightarrow |62-2\rangle \rightarrow |200\rangle$ [22]. The last state is rapidly absorbed, which we estimate

as $\Gamma_{200} \sim 10^{-11} \text{s}^{-1}$ for benchmark $\alpha = 0.124$ as relevant for OJ287 primary source. Thus, the boson cloud will be depleted in roughly $\sim 3 \times 10^3$ years once it transitions to the $|200\rangle$ state. However, during these resonances, the binary system is situated within the cloud. Full-state depletion would require $Z \gg 1$ in the preceding evolution steps. In realistic binaries the external energy loss is dominated by GW radiation and the cloud's dynamical friction, making γ large and $Z \lesssim 10^{-3}$. Consequently each resonance is highly nonadiabatic and transfers at most a few per cent of the state. This significantly suppresses the LZ parameter Z , resulting in an incomplete resonance transition.

In Fig. 6 we display the survival probability of the $|211\rangle$ and $|31-1\rangle$ levels after the chain of Bohr-type LZ crossings, as a function of (M_{BH}, μ) and for several characteristic BH binary mass ratios q . For equal masses ($q = 1$) the first crossing almost completely transfers state population from $|211\rangle$ to $|31-1\rangle$. However, $|31-1\rangle$ itself survives the subsequent crossing at the $> 99\%$ level, and the density profiles of the two states differ by less than an order of magnitude as shown on Fig. 2. When the companion is lighter, e.g., $q = 0.1$, more than 90% of the cloud remains in $|211\rangle$ over the entire parameter space, and in the residual region even a complete transfer to $|31-1\rangle$ would leave that state 100% intact considering $|31-1\rangle \rightarrow |62-2\rangle$ resonance. Considering the OJ287 mass ratio, $q \simeq 8 \times 10^{-3}$, the survival of $|211\rangle$ state exceeds 99% throughout its history. Hence, resonance depletion can be neglected in the analysis of OJ287.

Fine- and hyperfine-structure transitions require binary separations well above those reached by the OJ287 binary throughout its evolution history, for example, tracing back to redshift $z \simeq 10$, and therefore never occur. Backreaction of the few-percent state population that is shuffled between levels is negligible compared with the dominant dynamical-friction torque. Allowing for nonvanishing eccentricity or a $i \simeq 90^\circ$ inclination results in even greater survival fraction [24,73]. We therefore consider, for simplicity, a circular coplanar orbit in the main text.

Apart from the bound-bound resonance transition, there is also a resonance transition from a bound state to an unbound state, which is called ionization [24,51]. The backreaction of ionization was shown to be well described as dynamical friction [24]. Ionization can eject bosons and deplete the cloud, which, however, is subdominant compared with the accretion effect of the companion for a massive cloud such as the case we consider here [51]. A rough estimate of the timescale for accretion to consume the cloud is $\gtrsim 10^8$ yr, which is much larger than the duration of accretion experienced by the companion so far. We also estimate that tidal effects caused by the companion [74] result in a small tidal parameter $\epsilon \sim 10^{-8}$ for OJ287, indicating stability to them. Hence, dominant depletion channels are insignificant, and GA cloud can be considered as quasi-static, consistent with our analysis.

APPENDIX C: SMBH BINARY MERGER RATE

We can estimate the SMBH binary merger rate by mapping galaxy mergers to SMBH coalescences, considering each galaxy hosts a central BH. For primary SMBH mass by M_{BH} and mass ratio by q , and the primary galaxy stellar mass by M_g with merging galaxy mass ratio q_g , one has

$$\frac{d^3 n}{dz dq dM_{\text{BH}}} = \frac{d^3 n_g}{dz dq_g dM_g} \frac{dM_g}{dM_{\text{BH}}} \frac{dq_g}{dq}. \quad (\text{C1})$$

The galaxy merger rate is parametrized as [61,75]

$$\frac{d^3 n_g}{dz dq_g dM_g} = \frac{\Psi(M_g, z')}{M_g} \frac{\mathcal{P}(M_g, q_g, z')}{T_{g-g}(M_g, q_g, z')} \frac{dt}{dz'}, \quad (\text{C2})$$

where Ψ is the galaxy stellar mass function, $\mathcal{P}$ the galaxy pair fraction, and T_{g-g} the galaxy merger timescale. The cosmic time-redshift relation dt/dz in a flat Λ CDM cosmology is given by

$$\frac{dt}{dz} = \frac{1}{(1+z)H(z)}, \quad (\text{C3})$$

$$H(z) = H_0 \sqrt{\Omega_\Lambda + \Omega_m(1+z)^3}, \quad (\text{C4})$$

with Planck 2018 parameters $H_0 = 67.4 \text{ km s}^{-1} \text{ Mpc}^{-1}$, $\Omega_m = 0.315$, $\Omega_\Lambda = 0.685$ [76].

Adopting the Schechter form for galaxy stellar mass function, we parametrize the relevant quantities as

$$\Psi(M_g, z) = \Psi_0 \left(\frac{M_g}{M_\Psi} \right)^{s_\Psi} \exp \left(- \frac{M_g}{M_\Psi} \right), \quad (\text{C5})$$

$$\mathcal{P}(M_g, q_g, z) = \mathcal{P}_0 (1+z)^{\beta_{p0}}, \quad (\text{C6})$$

$$T_{g-g}(M_g, q_g, z) = T_0 (1+z)^{\beta_{t0}} q_g^{\gamma_{t0}}, \quad (\text{C7})$$

with

$$\log_{10} \left(\frac{\Psi_0}{\text{Mpc}^{-3}} \right) = \psi_0 + \psi_z z, \quad (\text{C8})$$

$$\log_{10} \left(\frac{M_\Psi}{M_\odot} \right) = m_{\psi_0} + m_{\psi_z} z, \quad (\text{C9})$$

$$s_\Psi = 1 + s_{\psi_0} + s_{\psi_z} z. \quad (\text{C10})$$

We can further relate M_g to the SMBH mass via the $M_{\text{BH}}-M_{\text{bulge}}$ relation [77]

$$\log_{10} \left(\frac{M_{\text{BH}}}{M_\odot} \right) = 8.7 + 1.1 \log_{10} \left(\frac{M_{\text{bulge}}}{10^{11} M_\odot} \right), \quad (\text{C11})$$

where the bulge mass can be expressed by the stellar mass of galaxy as [78,79]

TABLE II. Parameter values used for evaluating the SMBH binary merger rate.

Parameter	Value	Parameter	Value
ψ_0	-2.63	$\mathcal{P}_0$	0.033
ψ_z	-0.6	β_{p_0}	1
m_{ψ_0}	11.5	T_0 (Gyr)	0.5
m_{ψ_z}	0.11	β_{t_0}	-0.5
s_{ψ_0}	-1.21	γ_{t_0}	-1
s_{ψ_z}	-0.03	f_g	0.615

$$M_{\text{bulge}} = f_g M_g. \quad (\text{C12})$$

From Eqs. (C11) and (C12), the relation between the mass ratios q_g and q can be found as

$$q = q_g^{1.1}. \quad (\text{C13})$$

In Table II we display the numerical values of relevant parameters [61,62].

APPENDIX D: STATISTICAL ANALYSIS OF GW BACKGROUND FIT

To evaluate the performance of stochastic gravitational wave background in superradiant boson cloud scenario, we estimate the value of χ^2 in both superradiant boson cloud scenario and pure GW scenario. The χ^2 can be calculated as

$$\chi^2 = \sum_{i=1}^N \frac{(O_i - E_i)^2}{\sigma_i^2}, \quad (\text{D1})$$

where O_i is the i th observed characteristic strain at measured frequency, E_i is the corresponding expected value of characteristic strain in difference scenarios, and σ_i^2 is the variance of the i th distribution of characteristic strain. N is the total number of observed data points. Give the characteristic strain of PTA data as shown in Table III, we have the observed value of characteristic strain O_i and corresponding variance σ_i^2 . Since the value of upper variance σ_+^2 and lower variance σ_-^2 are different, we choose the mean value of upper and lower variance in χ^2 calculation as

$$\sigma^2 = \frac{\sigma_+^2 + \sigma_-^2}{2}. \quad (\text{D2})$$

Then we apply Eq. (21) to calculate E_i at measured frequency in O_i . It gives a χ^2 fit result from Fig. 5 that predicted strain in the GA cloud scenario gives $\chi^2 = 16.5$ (blue solid—all SMBHs with maximal spin), $\chi^2 = 17.5$ (blue dashed—conservative superradiance timescales), while predicted strain in GW scenario gives $\chi^2 = 18.8$ (black solid—general relativity prediction).

TABLE III. The data points of the PTA characteristic strain in NANOGrav [56], EPTA [64] and PPTA [65].

$f[\text{yr}^{-1}]$	$h_c/10^{-15}$	Ref	$f[\text{yr}^{-1}]$	$h_c/10^{-15}$	Ref
0.062	$6.5^{+4.5}_{-2.4}$	NANOGrav	0.055	$8.3^{+6.8}_{-3.4}$	PPTA
0.12	$7.9^{+3.0}_{-1.8}$		0.11	$9.6^{+4.2}_{-3.6}$	
0.19	$7.4^{+3.1}_{-2.0}$		0.17	$7.4^{+3.1}_{-1.8}$	
0.25	$6.4^{+3.3}_{-1.9}$		0.22	$6.3^{+4.5}_{-2.5}$	
0.31	$9.3^{+4.8}_{-4.2}$		0.28	$1.0^{+3.2}_{-0.8}$	
0.097	$8.0^{+4.0}_{-2.7}$	EPTA	0.33	$6.3^{+4.4}_{-4.0}$	
0.19	$9.6^{+2.9}_{-1.9}$		0.39	$3.7^{+3.7}_{-1.5}$	
0.29	$8.2^{+3.8}_{-8.2}$		0.44	$7.1^{+3.9}_{-2.6}$	
0.39	$11.0^{+4.5}_{-2.9}$		0.50	$1.6^{+3.2}_{-1.0}$	
0.48	$5.2^{+10.0}_{-5.5}$		0.55	$3.6^{+3.6}_{-2.09}$	
0.58	$0.6^{+4.8}_{-0.4}$				

APPENDIX E: BOSON SOLITONS AND GRAVITATIONAL WAVES

ULB DM can form compact solitonic cores around SMBHs. The resulting signatures imprinted on stochastic GW background allows restricting ULB masses in the range of 1.3×10^{-21} eV to 1.4×10^{-20} eV [66]. We calculate dynamical friction $\langle P_{\text{DF}} \rangle$ considering density

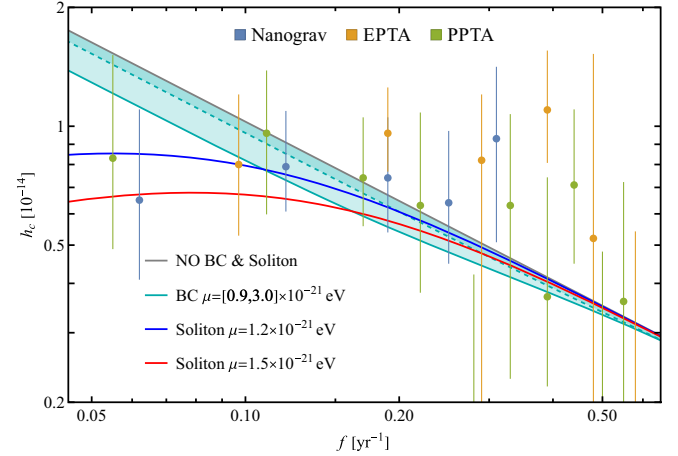


FIG. 7. The characteristic strain of GW background produced via SMBH binaries with different impacts from ULB soliton cores and superradiant GA boson clouds. We consider GA boson clouds in $|211\rangle$ state, with critical SMBH mass satisfying $\Gamma_{211} = 1/\tau_U$ for green solid line and $\Gamma_{211} = 170/\tau_U$ for green dashed line. The shaded region depicts impact from superradiant boson clouds on SGWB, considering ULB masses in the range of $(0.9-3) \times 10^{-21}$ eV. Impact of solitonic cores with ULB masses of $\mu = 1.2, 1.5 \times 10^{-21}$ eV is shown. The mass ratio ϵ between soliton core and SMBH is set as 10^{-3} . The observation data points come from NANOGrav [56], EPTA [64], and PPTA [65] are shown.

profile of solitons described as [66]

$$\rho_{\text{sol}} \simeq \rho_0 e^{-2a\mu r} \quad (\text{E1})$$

with

$$\rho_0 \simeq 5 \times 10^4 \varepsilon \left(\frac{M}{10^8 M_\odot} \right)^4 \left(\frac{\mu}{10^{-21} \text{ eV}} \right)^6 M_\odot \text{pc}^{-3}, \quad (\text{E2})$$

where ε is the mass ratio between solitonic core and SMBH.

In Fig. 7 we display the GW strain for stochastic background from SMBH solitonic cores for various

ULB DM masses considering population of SMBHs following Eq. (21) and SMBH mass distribution in the range of 10^5 – $10^{10} M_\odot$. We compare these results with dynamical friction effects from ULB GA boson clouds in [211] state considering SMBHs with mass greater than critical mass M_{crit} that satisfies condition $\Gamma_{211} = 1/\tau_U$. This leads to a much narrower SMBH population contribution to SGWB compared with the soliton case. The dynamical friction effects from ULB-DM solitonic cores and from superradiant clouds imprint qualitatively different low-frequency turnovers in the PTA spectrum, allowing these two scenarios to be distinguished.

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