

Superradiance constraints from GW231123

Andrea Caputo,^{1,2,*} Gabriele Franciolini^{1,†} and Samuel J. Witte^{3,4,5,‡}¹*Department of Theoretical Physics, CERN, Esplanade des Particules 1,
P.O. Box 1211, Geneva 23, Switzerland*²*Dipartimento di Fisica, “Sapienza” Università di Roma and Sezione INFN Roma1,
Piazzale Aldo Moro 5, 00185, Roma, Italy*³*Rudolf Peierls Centre for Theoretical Physics, University of Oxford, Oxford, United Kingdom*⁴*Deutsches Elektronen-Synchrotron DESY, Notkestraße 85, 22607 Hamburg, Germany*⁵*II. Institute of Theoretical Physics, Universität Hamburg,
22761, Hamburg, Germany*

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Gravitational wave observations have recently revealed with high significance, and high precision, the existence of $\mathcal{O}(100)M_\odot$ rapidly rotating black holes, allowing gravitational wave events to be used for the first time to probe unexplored axion parameter space using the phenomenon known as black hole superradiance. Here, we present new limits on axions using the binary black hole merger event GW231123, whose constituent black holes are among the fastest spinning observed with gravitational waves to date. We demonstrate that the most viable binary formation channels lead to conservative constraints on axion masses $\mu \sim [0.6 - 5] \times 10^{-13}$ eV and decay constants $f_\phi \gtrsim 10^{14}$ GeV, extending existing superradiance constraints derived using x-ray observations to yet lower axion masses.

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I. INTRODUCTION

The LIGO-Virgo-KAGRA (LVK) Collaboration has recently announced the detection of the merger of two constituent black holes (BHs) with masses $M_1 = 137^{+22}_{-17}M_\odot$ and $M_2 = 103^{+20}_{-52}M_\odot$ (at 90% credible intervals) [1]. From an astrophysical perspective, this event is remarkable, as the inferred masses of both constituents lie in the so-called “pair instability mass gap” [2–4] (i.e., the range of masses in which pair instabilities prevent the formation of BHs directly from stellar collapse) and represent the heaviest high-significance gravitational wave (GW) event observed to date [5,6]. This event is even more interesting from the perspective of a particle physicist, as there is a strong preference in the waveform that at least one, if not both, constituent BHs have extremely high dimensionless spin, $\tilde{a}_1 = 0.90^{+0.10}_{-0.19}$ and $\tilde{a}_2 = 0.80^{+0.20}_{-0.51}$, with the primary exhibiting the highest black hole spin

ever confidently measured through GWs. This unprecedented measurement is obtained thanks to the confident inference of a large effective precessing spin [7], a feature which appears robust against waveform systematics [1]. This is significant as it represents the first time GW data can be used to probe unexplored parameter space of new light fundamental bosonic particles via the phenomenon known as superradiance.

BH superradiance is a phenomenon in which low-energy bosonic fields experience an instability around Kerr BHs, causing them to rapidly transfer the rotational energy of the BH into high-occupation number quasibound states (see, e.g., [8] for a review). For a scalar field, this process, under optimal circumstances, can lead to an order 1 depletion of the BH spin on the timescale of hours [9–15], suggesting an inherent incompatibility between rapidly rotating BHs and the existence of light bosons [16–20].

To date, superradiance constraints have been derived using BH spin measurements that are inferred from x-ray observations [21–25]; this approach, however, can be highly sensitive to the properties of matter near the innermost stable circular orbit (where the modeling is the least well understood), the details of radiative transfer, and the thermodynamic properties of x-ray emitting regions—see, e.g., [21] for a detailed discussion of systematics. GW observations offer an alternative route to measuring high-spin BHs which is free from the modeling systematics that plague x-ray observations—the only difficulty, until now,

* Contact author: andrea.caputo@cern.ch† Contact author: gabriele.franciolini@cern.ch‡ Contact author: Samuel.Witte@physics.ox.ac.uk

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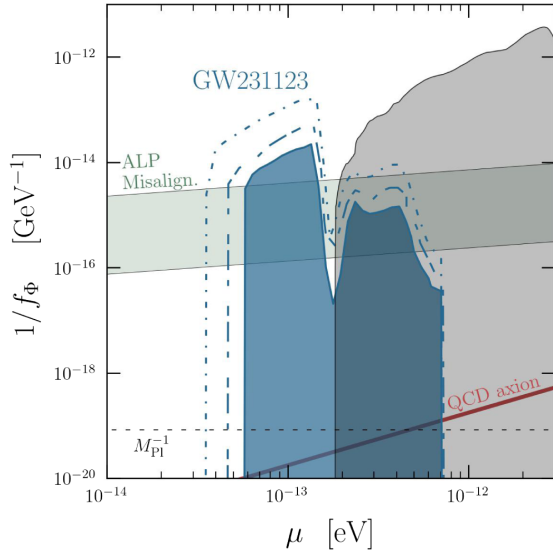


FIG. 1. Constraint on the axion decay constant f_Φ and mass μ derived from GW231123 (blue); result is obtained using NRSUR waveform model analysis, using $N_{\text{max}} = 5$, and with $\tau_{\text{max}} = 10^5$ (shaded), 10^6 (dashed), and 10^7 years (dot-dashed). Shown for comparison are superradiance constraints on solar-mass scale BHs [21], with spin inferences performed using x-ray data (black), the QCD axion line (red), and the parameter space where misalignment mechanism could produce the entirety of dark matter with $\mathcal{O}(1)$ initial field values (green). The inverse Planck scale M_{Pl}^{-1} is highlighted with a horizontal dashed line.

has been the lack of gravitationally observed high-spin BHs.¹

The aim of this paper is to demonstrate the power of GW231123, and more generally the power of forthcoming GW observations, in using spin measurements to constrain the existence of light bosonic particles. We focus specifically on the extent to which the inferred spin distributions of the constituent BHs in GW231123 can be used to probe axions, as these particles constitute the most well-motivated bosons at low energies (we do, however, derive a constraint on noninteracting spin-1 fields in the Supplemental Material [31], which includes Refs. [32–84]). Using the inferred properties of GW231123, we deduce the range of characteristic timescales over which superradiance could have spun down the BHs, within realistic assumptions motivated by the hierarchical formation scenario (see, e.g., [85] for a review). We then use the inferred posteriors on

¹There have been attempts to use radio observations of supermassive BHs [23]; however the large residual differences between analyses suggest uncertainties on the spin inference remain too large to be of use [26–29]. It is worth highlighting that during the review of this manuscript, LVK released an analysis of GW241011 [30], which yielded the most precise high-spin measurement to date; the mass of this event, however, is situated in the same range as BHs for which there exist x-ray spin measurements, and thus the derived constraints are subdominant to preexisting limits.

the spin of each constituent BH, along with motivated priors on their natal values, to simulate possible spin-down histories in the presence of an axion. The main result of this paper is the limit shown in Fig. 1, which demonstrates that GW231123 has excluded a new, and well-motivated, region of axion parameter space that is effectively inaccessible to conventional searches.² Needless to say, the constraints derived here rely on the inferred spins of the constituent BHs in GW231123; while the inference of high spins appears to be robust to variations in waveform modeling, the spin posteriors are still sensitive systematic uncertainties that can induce small shifts into the inferred limits (we note, however, that the sensitivity is primarily affected by the small-spin region of the posterior—see the Supplemental Material [31]). As waveform modeling improves in the future, so too will the robustness of these constraints; in addition, the approach and methodology used in the analysis here will undoubtedly prove invaluable as new detections of high-spin BHs emerge.

II. BLACK HOLE SUPERRADIANCE

Let us start with a brief review of BH superradiance. Here, we restrict our attention to scalar fields—generalizations to the case of higher spin fields can be found, e.g., in [8,11,12,14,90–98] (in the Supplemental Material [31], we also estimate constraints on noninteracting spin-1 fields). We focus specifically on the case of axions Φ , i.e., light pseudoscalars which have an intrinsic discrete shift symmetry $\Phi \rightarrow \Phi + 2\pi f_\Phi$ that protects the smallness of the mass, where f_Φ is the axion decay constant. The existence of the shift symmetry implies contributions to the axion potential are of the form $V(\Phi) = \mu^2 f_\Phi^2 [1 - \cos(\frac{\Phi}{f_\Phi})] + \dots$ (with higher order, shift-symmetry preserving terms, potentially being present), which for small field values reduces to $V(\Phi) \simeq \frac{1}{2}\mu^2 \Phi^2 - \frac{1}{4!}\lambda \Phi^4 + \dots$, with μ the axion mass, and $\lambda \equiv (\mu/f_\Phi)^2$. Thus, a natural expectation of axions is the existence of quartic self-interactions; despite a very small coupling constant, this interaction plays an important role in the evolution of superradiant systems [17,21,24,99–103].

The evolution of superradiant states can be determined by solving the equation of motion in a Kerr background. At leading order, one can neglect self-interactions and solve for the complex energy spectrum of quasibound states near the black hole; this energy spectrum is discrete, with eigenvalues $\omega = \omega_r + i\Gamma$ labeled by a set of quantum numbers $|\ell m\rangle$, see, e.g., [9,13,104]; these quantum

²GW evidence for highly spinning BHs has also been presented in [5,6,86–89]. In particular, GW190517 was also characterized by high spins. However, the component masses of that binary only probe regions of axion parameter space excluded by other constraints. Moreover, GW190517 had a modest network signal-to-noise ratio ($\text{SNR} \sim 10$), below the former detection threshold of 12 (although it does satisfy the new ranking statistic requirement of the false alarm rate $\text{FAR} < 0.25 \text{ yr}^{-1}$).

numbers are analogous to the principal n , orbital ℓ , and magnetic m quantum numbers of the hydrogen atom, obeying $n \geq \ell + 1$, $\ell \geq 0$, and $|m| \leq \ell$. The fastest growing state is the $|211\rangle$, which for a $\mathcal{O}(100)M_\odot$ BH has an e -fold growth timescale that can be as short as $\tau_{211} \equiv (2\Gamma_{211})^{-1} \sim \mathcal{O}(\text{hours})$. For a free particle, the growth of the state saturates when the angular frequency of the horizon Ω_H reaches a value $\Omega_H \equiv \tilde{a}/(2r_+) \sim \mu/m$,³ where $\tilde{a} \equiv a/(GM)$ is the dimensionless spin and $r_+ = GM(1 + \sqrt{1 - \tilde{a}^2})$ is the outer horizon. This requires roughly $\mathcal{O}(180)$ e -folds of growth, implying the spin of a rapidly rotating BH can be modified on timescales of days. In the calculations below, we compute the eigenfrequencies of each state and at each point during spin-down using the continued fraction method [10,105–107].

Self-interactions serve to couple the different eigenstates of the system, transferring energy efficiently between them.⁴ Rather than having decoupled growth of quasibound states, as in the case of the free particle solution, self-interactions induced coupled evolution of states in which the growth of any state can be quenched, saturating to a fixed occupation number, long before altering the spin of the BH [21,24,102,103]. The evolution of the superradiant system can be determined by following the energy injection and dissipation, which is described by the following set of equations:

$$\dot{\epsilon}_{n\ell m} = 2\Gamma_{n\ell m}\epsilon_{n\ell m} + \sum_{\mathcal{R}_{i,j,k}} \zeta_{i,j} \gamma_{i,j}^{k,L} \epsilon_i \epsilon_j \epsilon_k \quad (1)$$

$$\dot{\tilde{a}} = -\sum_{n\ell m} 2m\Gamma_{n\ell m}\epsilon_{n\ell m} \quad (2)$$

$$\frac{\dot{M}}{\mu GM^2} = -\sum_{n\ell m} 2\Gamma_{n\ell m}\epsilon_{n\ell m} + \sum_{(i,j,k)} \gamma_{i,j}^{k,BH} \epsilon_i \epsilon_j \epsilon_k, \quad (3)$$

where $\epsilon_{n\ell m} \equiv N_{n\ell m}/(GM^2)$ is the normalized occupation number $N_{n\ell m}$ of a state, and $\zeta_{i,j}$ is a degeneracy counting factor equal to 1 for $i \neq j$ and 2 for $i = j$. Here, we have introduced rate coefficients $\gamma_{i,j}^{k,L}$, where the indices (i, j, k) denote quasibound states (i.e., $i = |n\ell m\rangle$, $j = |n'\ell'm'\rangle$, etc), and “L” represents energy loss, either to infinity ∞ (occurring when $\omega_i + \omega_j - \omega_k > \mu$), or into bound states which dissipate their energy into the BH (occurring when $\omega_i + \omega_j - \omega_k \leq \mu$). The collective list of nonzero scattering rates relevant for $n \leq 5$ is presented in Tables 1 and 2 of [21]. Note that the axion mass implicitly enters in both the

eigenvalues $\Gamma_{n\ell m}$ as well as the scattering rates $\gamma_{i,j}^{k,L}$, while the decay constant enters only through the scattering rates.

In general, one should include all states that obtain non-negligible occupation numbers on the timescales of interest; in practice, however, this is currently unfeasible, with the latest analyses only including states with $n \leq 5$ [21]. Here, we perform our fiducial analysis with all states $n \leq 5$, and illustrate the impact of taking instead $n \leq 3$ in the Supplemental Material [31]. All rate coefficients are computed using the procedure of [21].

In solving Eqs. (1)–(3), there is an important subtlety which arises from the fact that we observe the spin and mass close to the merger, and these values may differ sizably from those at birth. For this reason, one typically evolves the superradiant system on timescales $t \lesssim \tau_{\text{max}}$, where τ_{max} is chosen to be sufficiently short that astrophysical processes cannot appreciably alter the mass or spin of the system; for isolated BHs, this is often taken to be the Salpeter timescale, $\tau_{\text{Salp}} \sim 4.5 \times 10^7$ yr; however we argue in the following section that for GW231123 the relevant timescale may be shorter, and thus we adopt a more conservative timescale in our analyses. In general, one could attempt to model the evolution on longer timescales, including, e.g., accretion, multiple mergers, etc. [118], but this requires complicated astrophysical modeling which is beyond the scope of this work.

III. DETERMINATION OF MERGER TIMESCALES

As mentioned in the preceding section, testing superradiance typically requires defining the characteristic timescale over which the properties of the black hole can be assumed to remain constant. For isolated black holes, this is often taken to be the Salpeter timescale. For second or third generational hierarchical mergers, the relevant timescale is instead determined by the time between successive mergers, τ_{merger} , which can be significantly shorter than τ_{Salp} . Since the large masses and spins of the BHs involved in the GW231123 event strongly suggest that one or both of the binary components may be the product of a previous binary BH merger [1,119,120], a scenario known as a hierarchical merger [121–134] (see, e.g., [85] for a review), we endeavor to determine the minimal value of τ_{merger} consistent with the properties of GW231123. In the Supplemental Material [31], we further discuss more exotic scenarios, showing the constraint in Fig. 1 remains valid.

Assuming a hierarchical origin, one can infer important information about the general properties of the local environment. For instance, the presence of high-spin BHs in GW231123 is consistent with significant anisotropic GW emission, which implies large recoil velocities for the remnants of first-generation mergers which constitute the progenitors of GW231123—typically on the order of several hundred kilometers per second [135]. If this “birth kick” exceeds the escape velocity of the host environment, further

³We adopt natural units throughout this work.

⁴Generally speaking, other interactions could play a role in quenching superradiant growth (see, e.g., [24,108–117]); however the large occupation numbers N of quasibound states stimulate the $2 \leftrightarrow 2$ scattering process, causing the rate of energy dissipation to scale as $\propto N^3$ ensuring it dominates the energy dissipation.

mergers would be prevented. This suggests that GW231123 likely originated in a deep potential well, e.g., in a nuclear star cluster (NSC) [123,133,136–140] or an active galactic nucleus (AGN) [141–152], where escape velocities can reach $v_{\text{esc}} \sim 100\text{--}500$ km/s or higher. In contrast, typical escape velocities from young stellar clusters and globular clusters are significantly lower, $v_{\text{esc}} \lesssim 50$ km/s [153–155], making it unlikely that the remnants of earlier mergers could have been retained in such environments.

In general, one can write the rate density of BH-BH mergers as

$$\mathcal{R} = n_{\text{env}} N_{\text{BH}}(M_{\text{BH}}) \Gamma_{\text{merger}}, \quad (4)$$

where n_{env} is the average number density in the Universe of the environment under consideration, N_{BH} is the number of BHs within such an environment (where we included explicitly the dependence on the BH mass), and Γ_{merger} is the merger rate per BH. Such a rate can be decomposed into two contributions, $\Gamma_{\text{merger}} \sim \frac{1}{\tau_{\text{pair}} + \tau_{\text{coal}}}$, where τ_{pair} is the pairing timescale for two BHs, and τ_{coal} is the time to coalescence once the binary has formed. In what follows, we will assume the former dominates the merger timescales and that the coalescing time is negligible. This is particularly justified for GW231123 which is expected to form from dynamical processes involving large velocity BHs. In fact, one can verify that for the parameters of interest, namely $v \sim v_{\text{recoil}} \gtrsim 100$ km/s, $\tau_{\text{pair}} \gg \tau_{\text{coal}}$, where v denotes the relative velocity of a BH pair and v_{recoil} represents the “birth-kick” recoil velocity of the remnants from first-generation mergers. Such a hierarchy is naturally expected, as both the binary formation rate (scaling as $\tau_{\text{pair}} \propto v^{11/7}$ for dynamical capture [148,156], and $\tau_{\text{pair}} \propto v^9$ for three-body formation [157,158]) and the coalescing time (scaling as $\tau_{\text{coal}} \propto v^{-36/7}$ for dynamical capture, and $\tau_{\text{coal}} \propto v^{-8}$ for three-body formation, assuming Peter’s formula [159,160]) scale strongly with v . A similar argument applies to other channels, such as, e.g., binary-single interactions. We will proceed under this assumption, which allows us to treat superradiance growth around both individual isolated BHs (see Supplemental Material [31] for a discussion on the impact of the binary), before they pair in the binary and merge soon after.

The rate density in Eq. (4) should then be compared with the one inferred by the LVK Collaboration for GW231123, $\mathcal{R}_{\text{obs}} = 0.08^{+0.19}_{-0.07} \text{Gpc}^{-3} \text{yr}^{-1}$. From such a comparison, one can infer the characteristic pairing time,

$$\tau_{\text{pair}} \simeq \tau_{\text{merger}} \equiv \Gamma_{\text{merger}}^{-1} = \frac{n_{\text{env}} N_{\text{BH}}(M_{\text{BH}})}{\mathcal{R}_{\text{obs}}}. \quad (5)$$

This estimate is fully agnostic with respect to the physical mechanisms responsible for binary formation and coalescence. In order to provide realistic numbers, we will evaluate

Eq. (5) considering at least one of the BHs to be a remnant, and treating separately the possibility that this event occurred in an AGN or NSC.

A. Active galactic nuclei

In AGN, gas-rich accretion disks surround supermassive BHs, and stellar-mass BHs can become embedded and experience migration due to torques from the surrounding gas [149]. A particularly relevant feature of AGN disks is the presence of *migration traps*—locations where the net torque acting on an object vanishes, halting its inward or outward drift [161]. These traps naturally concentrate compact objects, enhancing the probability of dynamical interactions and subsequent mergers, making AGN disks ideal for multiple merger generations. As such, they offer a promising, natural scenario for explaining the high-spin, hierarchical nature of events like GW231123.

In this environment, the merger timescale in Eq. (5) is

$$\tau_{\text{merger}}^{\text{AGN}} = \text{Myr} \left(\frac{n_{\text{env}}^{\text{AGN}}}{10^{-5} \text{Mpc}^{-3}} \right) \left(\frac{N_{\text{BH}}}{10} \right) \times \left(\frac{0.1 \text{Gpc}^{-3} \text{yr}^{-1}}{\mathcal{R}_{\text{obs}}} \right), \quad (6)$$

where for the AGN number density at low redshift ($z \lesssim 0.5$), we adopted a typical, yet conservative, value inferred from multiwavelength luminosity functions (see, e.g., Fig. 7 in Ref. [162]). This number density only includes contributions from AGN with luminosities within the range where migration traps are expected to form and persist [130]. The number of massive, second-generation BHs in the trap is instead estimated following Ref. [163] (see their Fig. 9), where the authors simulate first-generation BH mergers embedded in the AGN disk and track their migration, binary pairing, and mergers. We stress that if the number of BHs is larger than expected, for a fixed observed rate, this would imply a larger merger timescale, and therefore one could probe a larger region of the axion parameter space. We also note that the observed rate, $\mathcal{R}_{\text{obs}}$, ranges from 0.01 to $0.3 \text{Gpc}^{-3} \text{yr}^{-1}$ at 90% confidence level, implying that τ_{pair} could be smaller by a factor of a few, or alternatively up to ten times longer. Overall, adopting

$$\tau_{\text{merger}}^{\text{AGN}} \sim 10^{-1} - 10 \text{ Myr} \quad (7)$$

should abundantly encompass all uncertainties.

B. Nuclear star clusters

NSCs [164] represent an alternative dynamical environment in which compact-object binaries can form and merge multiple times. Compared to AGN disks, NSCs are gas-poor systems, but they can nonetheless sustain high merger rates due to their large stellar densities and deep gravitational potentials. The dominant binary formation mechanism in

NSCs depends on whether or not a central supermassive BH (SMBH) is present. In NSCs that do host an SMBH, the short relaxation time facilitates the development of a steep density cusp of stellar-mass BHs around the central object, enhancing the probability of binary formation via gravitational-wave capture [148]. In contrast, NSCs without a SMBH behave more like high-mass globular clusters, where binaries form and harden primarily through three-body encounters and binary single interactions [122]. In either case, we can write an expression for the merger rate analogous to that in Eq. (4).

The number density of NSCs in the local Universe, relevant for events like GW231123, whose redshift is $z = 0.39^{+0.27}_{-0.24}$, depends on the fraction of galaxies that host NSCs and on the galaxy stellar mass function. The former peaks at approximately 90% for galaxies with a stellar mass around $\sim 10^9 M_\odot$, and decreases toward both higher and lower masses. Nonetheless, nucleation fractions remain at the level of $\sim 10\%$ even for galaxies as small as $M_* \sim 5 \times 10^5 M_\odot$ and as massive as $M_* \sim 5 \times 10^{11} M_\odot$ (see Fig. 3 of Ref. [165]). Each galaxy can host at most one NSC, located at its photometric and dynamical center (unlike the case for globular clusters). This allows us to estimate $n_{\text{env}}^{\text{NSC}} \sim n_{\text{gal}} f_{\text{NSC}}$. Using results from galaxy stellar mass function studies [166–168], and adopting a conservative value of $f_{\text{NSC}} \sim 10\%$, we obtain (see also [136]) $n_{\text{env}}^{\text{NSC}} \sim 10^{-3} \text{ Mpc}^{-3}$.

As in the case of AGN, determining $N_{\text{BH}}(M_{\text{BH}})$ is less straightforward, but we can expect it to be in the same ballpark. In fact, assuming standard stellar initial mass functions [169], one expects 0.1–1% of the NSC mass to be in the form of seed BHs that may form binaries and merge [170,171]. Typical NSC have masses $\sim 10^5\text{--}10^8 M_\odot$, which implies roughly $M_{\text{BH,tot}} \sim 10^3\text{--}10^4 M_\odot$. Of these BHs, we can expect 10–50% of them to be in binaries [172]. The majority of these BH will on average have masses $\sim 10\text{--}20 M_\odot$, leaving us with $\sim \mathcal{O}(10\text{--}100)$ BHs which will then undergo repeated mergers. We can thus write

$$\tau_{\text{merger}}^{\text{NSC}} = 10 \text{ Myr} N_{\text{BH}} \left(\frac{n_{\text{env}}^{\text{NSC}}}{10^{-3} \text{ Mpc}^{-3}} \right) \left(\frac{0.1 \text{ Gpc}^{-3} \text{ yr}^{-1}}{\mathcal{R}_{\text{obs}}} \right), \quad (8)$$

implying

$$\tau_{\text{merger}}^{\text{NSC}} \sim 1 - 100 \text{ Myr} \quad (9)$$

should bracket all uncertainties.

IV. ANALYSIS AND CONCLUSIONS

Here, we begin by providing a heuristic description of our statistical framework that allows one to derive constraints on the axion parameter space; a statistically rigorous derivation

of the relevant quantities from the LVK posterior is also included for completeness.

First, we fix the axion mass; this is useful as our posterior will otherwise be unbounded in three directions (and thus risks having a high sensitivity to the prior). We adopt a log-flat prior on f_ϕ over the range $[10^{11}, 10^{20}] \text{ GeV}$. We focus on understanding the evolution of the BHs from a time $t = t_0 - \tau_{\text{max}}$, to the time of merger t_0 , as uncertainties in the merger history do not allow us to infer the properties of the BHs on longer timescales. In our analysis, we vary τ_{max} to be consistent with the shortest characteristic timescales identified in the viable formation channels associated with NSCs and AGNs, with our fiducial limits taken to be the minimum of these, namely 10^5 years. Since the change in mass of each BH due to superradiance is small relative to the mass uncertainty, one can account for the mass measurement by adopting mass priors using the LVK joint prior distribution $p(M_1, M_2)$. The initial spins at a time $t = t_0 - \tau_{\text{max}}$ are, of course, unknown—while astrophysical processes cannot alter the spins on these timescales, superradiance can. Since superradiance only serves to extract spin, we adopt spin priors which contain nonzero weight only above the inferred values of the spins at merger. More specifically, given a sample i of the $X = 1, 2$ BH (with “1” corresponding to the heavier BH), we compute the conditional distribution $p(\tilde{a}_X | M_{1,i}, M_{2,i})$; we then determine the mean ζ_X and standard deviation β_X in the conditional probability $p(\tilde{a}_X | M_{1,i}, M_{2,i})$ (which we infer from the LVK posterior), and adopt a prior that is half flat, for $\tilde{a}_{X,i} > \zeta_X$, and half Gaussian, for the initial spin $\tilde{a}_{X,i} \leq \zeta_X$. The half-Gaussian distribution ensures an appropriate sampling of the initial spin distribution in the event that superradiance does not spin down the black hole, while the half-flat component of the prior ensures that initial high-spin black holes which have been spun down via superradiance to $\tilde{a} = \langle \tilde{a}_X \rangle$ are perfectly compatible with the observed data.⁵ Equations (1)–(3) are then evolved from $t = t_0 - \tau_{\text{max}}$ to $t = t_0$, to determine the spins that would have been observed by LVK $\tilde{a}_{1,2}$ at merger. In order to evaluate the compatibility of the predicted spins with those observed, we construct a likelihood using two-dimensional kernel density estimation on the set of conditioned posterior samples $p(\tilde{a}_1, \tilde{a}_2 | M_{1,i}, M_{2,i})$ from the LVK analysis. The interested reader can find a more detailed derivation of the statistical procedure in the Supplemental Material [31].

We run between five and ten Monte Carlo chains, reconstruct the one-dimensional marginalized posterior on f_ϕ from the Monte Carlo samples, and extract the

⁵Within the hierarchical scenario, the prior distribution of spin for the multigeneration mergers can be derived from specific assumptions on the progenitors and the environment; see, e.g., [119,173–175]. In most cases, the distribution is rather peaked around $\tilde{a} \sim 0.7$, as obtained for a remnant of mergers of spinless BHs. However, this is reasonably broad assuming spinning progenitors and/or higher-generation BHs, supporting our choice.

90% lower limit. This procedure is repeated over many axion masses, and the collection of these points produces the limit shown in Fig. 1. Note the multibump feature is driven primarily by the contributions of the heavier (left) and lighter (right) BHs (while multibumped features do also arise from the presence of different superradiant states for single BHs, the $|322\rangle$ state of the heavy BH is subdominant to the $|211\rangle$ state of the light BH at higher axion masses). For reference, we also show in Fig. 1 superradiance constraints (gray) derived using x-ray observations [21], the values of $f_\phi - \mu$ corresponding to the QCD axion (red), and the parameter space where the misalignment mechanism can most naturally generate the dark matter abundance (green). Note that the limit derived in Fig. 1 contains sensitivity to *both* black holes (see the Supplemental Material [31]). The various features seen in the limit arise from the interplay of having BHs of different masses, and looking at the spin-down induced by the $m = 1$ and $m = 2$ states.

This analysis demonstrates that GW observations of high-spin BHs can be used to set competitive limits on axions (and, more generally, light bosonic fields) in unexplored regions of the parameter space. The dominant uncertainty entering this analysis comes from the unknown timescale between successive mergers of the constituent BHs—this is in contrast to approaches using x-ray observations, where the dominant systematic arguably comes from the modeling of the accreting material near the horizon. Using a data-driven approach, we have argued that this uncertainty is not prohibitive, and that GW231123 can be used to derive meaningful constraints on feebly coupled axions (see [176] for a similar conclusion).

The upcoming high-sensitivity LVK O5 run [177] promises to significantly enhance discovery potential for similar events. Thanks to an expected signal-to-noise ratio improvement of $\mathcal{O}(3)$ [178], the observable volume will expand by nearly $\mathcal{O}(20)$. The detection of additional events of this kind in future LVK runs would be accompanied by narrower uncertainties on event parameters and will allow for improved constraints on the underlying spin distribution, along with other source population properties, thereby strengthening the bounds discussed in this work.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [179].

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