Black-Hole Bombs and Photon-Mass Bounds

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Phys. Rev. Lett. **109**, 131102 — Published 27 September 2012

DOI: 10.1103/PhysRevLett.109.131102

Black hole bombs and photon mass bounds

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Generic extensions of the standard model predict the existence of ultralight bosonic degrees of freedom. Several ongoing experiments are aimed at detecting these particles or constraining their mass range. Here we show that massive vector fields around rotating black holes can give rise to a strong superradiant instability which extracts angular momentum from the hole. The observation of supermassive spinning black holes imposes limits on this mechanism. We show that current supermassive black hole spin estimates provide the tightest upper limits on the mass of the photon ($m_v \lesssim 4\times10^{-20}$ eV according to our most conservative estimate), and that spin measurements for the largest known supermassive black holes could further lower this bound to $m_v \lesssim 10^{-22}$ eV. Our analysis relies on a novel framework to study perturbations of rotating Kerr black holes in the slow-rotation regime, that we developed up to second order in rotation, and that can be extended to other spacetime metrics and other theories.

PACS numbers: 04.40.Dg, 04.62.+v, 95.30.Sf

Introduction. The properties of matter making up our universe are mostly unknown. Strong evidence (e.g. from galactic rotation curves and from gravitational lensing) points to the existence of elusive, weakly-interacting matter as the most abundant element in the universe. An interesting possibility is the existence of ultralight bosonic degrees of freedom, such as those appearing in the "string axiverse" scenario [1, 2], or of massive hidden $U(1)$ vector fields, that are also a generic feature of extensions of the standard model [3–6].

Massive fields around rotating black holes (BHs) can trigger a superradiant instability, the so-called "black hole bomb" [7]. This instability is well understood in the case of massive scalar fields [8–15]: it requires the existence of negative energy states in a region around the BH known as the ergoregion. The instability is regulated by the dimensionless parameter $M\mu$ (from now on we set $G = c = 1$), where M is the BH mass and $m_s = \mu\hbar$ is the scalar field mass, and it is most effective when $M\mu \sim 1$ and for maximally spinning BHs. For a solar mass BH and a field of mass $m_s \sim 1$ eV the parameter $M\mu \sim 10^{10}$, and therefore in many cases of astrophysical interest the instability timescale is larger than the age of the universe. Superradiant instabilities strong enough to be observationally relevant ($M\mu \sim 1$) can occur either for light primordial BHs that may have been produced in the early universe [16–18] or for ultralight exotic particles found in some extensions of the standard model [1, 2]. In the string axiverse scenario, massive scalar fields with 10^{-33} eV $< m_s < 10^{-18}$ eV could play a key role in cosmological models. Superradiant instabilities may allow us to probe the existence of ultralight bosonic fields by producing gaps in the mass-spin BH Regge spectrum [1, 2], by modifying the inspiral dynamics of compact binaries [15, 19, 20] or by inducing a "bosenova", i.e. a collapse of the axion cloud [21–23].

The curved spacetime dynamics of massive vector fields has only been studied in nonrotating backgrounds [24, 25]. While superradiant instabilities are expected to occur also for massive vector fields, quantitative investigations have been hampered by our inability to fully understand the (massive vector) Proca equation

$$\nabla_\sigma F^{\sigma\rho} - \mu^2 A^\rho = 0\,, \tag{1}$$

where A_μ is the vector potential, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ and $m_v = \mu\hbar$ is the mass of the vector field. Note that the Lorenz condition $\nabla_\mu A^\mu = 0$ is automatically satisfied and the Proca field A_μ propagates 3 degrees of freedom [24]. In a nutshell, the problem is that Eq. (1) does not seem to be separable in the Kerr background.

Framework. The Proca perturbation problem in the Kerr metric becomes tractable if we work in the slow-rotation approximation. Let us focus on the Kerr metric in Boyer-Lindquist coordinates

$$ds^2_{\rm Kerr} = -\left(1-\frac{2Mr}{\Sigma}\right)dt^2 + \frac{\Sigma}{\Delta}dr^2 - \frac{4rM^2}{\Sigma}\tilde{a}\sin^2\vartheta d\varphi dt + \Sigma d\vartheta^2 + \left[(r^2+M^2\tilde{a}^2)\sin^2\vartheta + \frac{2rM^3}{\Sigma}\tilde{a}^2\sin^4\vartheta\right]d\varphi^2\,,$$

where $\Sigma = r^2 + M^2\tilde{a}^2\cos^2\vartheta$, $\Delta = (r-r_+)(r-r_-)$, $r_\pm = M(1\pm\sqrt{1-\tilde{a}^2})$ and M and $J = M^2\tilde{a}$ are the mass

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and the angular momentum of the BH, respectively. In what follows, we shall expand the metric and all other quantities of interest to second order in $\tilde{a}$. The procedure to separate the linearized perturbation equations to first order in $\tilde{a}$ was first proposed by Kojima in the context of stellar perturbations [26–28], but it can be generalized to *any* order in $\tilde{a}$ and to generic (scalar, vector, tensor, etc.) perturbations of stationary and axisymmetric spacetimes. The details of the procedure will appear elsewhere [29]; here we only present the main results.

In the slow-rotation limit the perturbation equations, expanded in spherical harmonics and Fourier transformed in time, yield a coupled system of ODEs. In the case of a spherically symmetric background, perturbations with different harmonic indices (ℓ, m), as well as perturbations with opposite parity, are decoupled. In a rotating, axially symmetric background, perturbations with different values of the azimuthal number m are still decoupled, but perturbations with different values of ℓ are not. However, in the limit of slow rotation there is a Laporte-like "selection rule" [30]: at first order in $\tilde{a}$, perturbations with a given value of ℓ are only coupled with those with $\ell \pm 1$ and *opposite* parity, similarly to the case of rotating stars. At second order, perturbations with a given value of ℓ are also coupled with those with $\ell \pm 2$ and *same* parity, and so on.

In general, the perturbation equations can always be written in the form [29]

$$\begin{aligned}
0 = &\mathcal{A}_\ell + \tilde{a}m\bar{\mathcal{A}}_\ell + \tilde{a}^2\hat{\mathcal{A}}_\ell \\
&+ \tilde{a}(\mathcal{Q}_\ell\tilde{\mathcal{P}}_{\ell-1} + \mathcal{Q}_{\ell+1}\tilde{\mathcal{P}}_{\ell+1}) \\
&+ \tilde{a}^2\left[\mathcal{Q}_{\ell-1}\mathcal{Q}_\ell\breve{\mathcal{A}}_{\ell-2} + \mathcal{Q}_{\ell+2}\mathcal{Q}_{\ell+1}\breve{\mathcal{A}}_{\ell+2}\right] + \mathcal{O}(\tilde{a}^3)\,, \quad (2)\\
0 = &\mathcal{P}_\ell + \tilde{a}m\bar{\mathcal{P}}_\ell + \tilde{a}^2\hat{\mathcal{P}}_\ell \\
&+ \tilde{a}(\mathcal{Q}_\ell\tilde{\mathcal{A}}_{\ell-1} + \mathcal{Q}_{\ell+1}\tilde{\mathcal{A}}_{\ell+1}) \\
&+ \tilde{a}^2\left[\mathcal{Q}_{\ell-1}\mathcal{Q}_\ell\breve{\mathcal{P}}_{\ell-2} + \mathcal{Q}_{\ell+2}\mathcal{Q}_{\ell+1}\breve{\mathcal{P}}_{\ell+2}\right] + \mathcal{O}(\tilde{a}^3)\,, \quad (3)
\end{aligned}$$

where $\mathcal{Q}_\ell = \sqrt{\frac{\ell^2-m^2}{4\ell^2-1}}$ and the coefficients $\mathcal{A}_\ell$ and $\mathcal{P}_\ell$ (with various superscripts) are *linear* combinations of axial and polar perturbation variables, respectively.

The general method can be specialized to the Proca equation (1). We expand the vector potential as [24]

$$\delta A_\mu(t,r,\vartheta,\varphi) = \sum_{\ell,m}\left[\begin{array}{c} 0 \\ 0 \\ u^\ell_{(4)}\mathbf{S}^\ell_b/\Lambda \end{array}\right] + \left[\begin{array}{c} u^\ell_{(1)}Y^\ell/r \\ u^\ell_{(2)}Y^\ell/(rf) \\ u^\ell_{(3)}\mathbf{Y}^\ell_b/\Lambda \end{array}\right]\,,$$

where $b \equiv (\vartheta, \varphi)$, $\Lambda = \ell(\ell+1)$, $f = \Delta/(r^2 + \tilde{a}^2M^2)$, Y^ℓ are scalar spherical harmonics, $\mathbf{Y}^\ell_b \equiv (Y^\ell_\vartheta, Y^\ell_\varphi)$ and $\mathbf{S}^\ell_b \equiv (S^\ell_\vartheta, S^\ell_\varphi)$ are vector spherical harmonics [29], whereas $u^\ell_{(i)} = u^\ell_{(i)}(r,t)$ $(i = 1,2,3)$ and $u^\ell_{(4)} = u^\ell_{(4)}(r,t)$ are polar and axial perturbations, respectively. Separating the angular variables, we find that Proca perturbations in the slow-rotation limit, up to second order, are described by two sets of equations [29]:

$$\mathcal{D}_\mathbf{A}\mathbf{\Psi_A}{}^\ell + \mathbf{V_A}\mathbf{\Psi_A}{}^\ell = 0\,, \tag{4}$$

$$\mathcal{D}_\mathbf{P}\mathbf{\Psi_P}{}^\ell + \mathbf{V_P}\mathbf{\Psi_P}{}^\ell = 0\,, \tag{5}$$

where $\mathcal{D}_{\mathbf{A},\mathbf{P}}$ are second order differential operators, $\mathbf{V}_{\mathbf{A},\mathbf{P}}$ are matrices, $\mathbf{\Psi_A}{}^\ell = (u^\ell_{(4)}, u^{\ell\pm1}_{(2)}, u^{\ell\pm1}_{(3)}, u^{\ell\pm2}_{(4)})$ and $\mathbf{\Psi_P}{}^\ell = (u^\ell_{(2)}, u^\ell_{(3)}, u^{\ell\pm1}_{(4)}, u^{\ell\pm2}_{(2)}, u^{\ell\pm2}_{(3)})$. The function $u^\ell_{(1)}$ can be obtained from the Lorenz condition once the three dynamical degrees of freedom are known [29]. When $\tilde{a} = 0$, the equations above reduce to Proca perturbations of a Schwarzschild BH [24]. However, rotation introduces mixing between perturbations of different parity and different multipolar indices.

Numerical Results. Once suitable boundary conditions and a time dependence of the form $e^{-\mathrm{i}\omega t}$ are imposed, Eqs. (4)–(5) form an eigenvalue problem for the complex frequency $\omega = \omega_R + \mathrm{i}\omega_I$. Physically motivated boundary conditions correspond to either quasinormal modes (perturbations having ingoing-wave conditions at the horizon and outgoing-wave conditions at infinity [31]) or bound states (perturbations that are spatially localized within the vicinity of the BH and decay exponentially at infinity). Here we focus on bound modes. By analogy with the scalar field case we would expect these modes to become superradiantly unstable for $\omega_R < m\Omega_H$ [10], where $\Omega_H = \tilde{a}/(2r_+)$. Our analysis shows, for the first time, that massive vector fields do indeed become unstable in this regime.

The bound state modes of the system (4)–(5) can be found by standard numerical methods [29]. When $m > 0$ we find that, within numerical errors, the imaginary part of the modes has a zero crossing when

$$\omega_R = m\Omega_H \sim m\frac{\tilde{a}}{4M} + \mathcal{O}(\tilde{a}^3)\,, \tag{6}$$

which corresponds to the onset of the superradiant regime. When $\omega_R < m\Omega_H$ the modes are unstable and $\tau = \omega_I^{-1}$ is the instability growth timescale. Note that, although the superradiant condition (6) appears as a first order effect, in fact $\omega M \sim \tilde{a}$ at the onset of superradiance. The field equations contain terms proportional to $\tilde{a}\omega$ and to ω^2, so a second-order expansion is needed for a self-consistent study of the unstable regime [29].

When $M\mu \lesssim 0.1$ our data for the fundamental modes are consistent with a hydrogenic spectrum, $\omega_R \sim \mu$ and

$$M\omega_I \sim \gamma_{S\ell}\,(\tilde{a}m - 2r_+\mu)\,(M\mu)^{4\ell+5+2S}\,, \tag{7}$$

where $\gamma_{S\ell}$ is a coefficient that depends on ℓ and on the "polarization" index S, with $S = 0$ for axial modes and $S = \pm1$ for two classes of polar modes [24]. In the axial case the numerical results are also supported by an analytical formula, that can be found by applying Starobinski's method of matching asymptotics [10, 32] to Eq. (4) at first order. A detailed calculation [29] yields precisely Eq. (7) with $S = 0$ and $\gamma_{01} = 1/12$. This is consistent with our numerical data, that yield $\gamma_{01} \approx 0.09 \pm 0.03$

(here and in the following, numerical errors are estimated by comparing the results at first and second order and by taking the maximum deviation between the fit and the data). The instability timescale in the axial Proca case is four times shorter than in the scalar case. Similar results can be derived also for axial modes with $\ell > 1$ and for the overtones.

To our knowledge, this is the first estimate of the instability timescale of massive vector fields around spinning BHs. Although Eq. (7) is strictly valid only when $\tilde{a} \ll 1$ and $M\mu \ll 1$, in the case of massive scalar fields it provides estimates in good agreement with exact results [12, 13] up to $\tilde{a} \lesssim 0.99$: for example, Eq. (7) overestimates the exact result by only 3% when $\tilde{a} = 0.7$, and by less than 70% when $\tilde{a} = 0.99$ [29].

In the massive vector case, the good agreement between the first- and second-order calculations suggests that the slow-rotation expansion can be trusted even for moderately large spins, $\tilde{a} \lesssim 0.7$. Therefore it is reasonable to expect that extrapolations of Eq. (7) from the slow-rotation limit should at least provide the correct order of magnitude (and possibly a reliable quantitative estimate) of the instability timescale far from extremality. If we extrapolate Eq. (7) to $\tilde{a} \to 1$, we expect to overestimate the instability by about one order of magnitude [29].

As we shall discuss below, astrophysical bounds on vector field masses (scaling as $\gamma_{S\ell}^{-1/(4\ell+5+2S)}$ [29]) depend very mildly on uncertainties in the $\gamma_{S\ell}$ coefficients. However they are sensitive to the scaling with μ in the $M\mu \ll 1$ regime, and it is crucial to obtain reliable results in this limit. Unfortunately the calculation of unstable modes when $M\mu \lesssim 0.02$ is challenging due to numerical inaccuracies, but the consistency between our numerical data in the axial case and the analytical formula (which is valid when $M\mu \ll 1$) is reassuring. The equations governing polar modes are much more complex, and we could not find analytical results supporting the fit (7) when $S \neq 0$ (cf. [24]). Our data for $S = -1$ polar modes at the onset of the $\ell = m = 1$ instability are consistent with Eq. (7) with $\gamma_{-11} \approx 20 \pm 10$, but some experimentation showed that different fitting functions can provide even better fits. Due to these uncertainties in the fit (7) for the polar case, in the following we shall discuss the consequences of our results mainly for axial modes, for which we derived the instability timescale also *analytically* [29]. This choice is very conservative because, according to Eq. (7), polar modes with $S = -1$ exhibit the strongest instability. Indeed, for fixed values of $\tilde{a}$ and μ the instability of polar modes with $S = -1$ is typically two or three orders of magnitude stronger than in the axial case [29].

Astrophysical bounds on the photon mass. Our results, together with reliable supermassive BH spin measurements, can be used to impose stringent constraints on the allowed mass range of massive vector fields. These bounds follow from the requirement that astrophysical spinning BHs should be stable, in the sense that the instability timescale τ should be larger than some observational threshold. For isolated BHs we can take the observational threshold to be the age of the Universe, $\tau_{\rm Hubble} = 1.38 \times 10^{10}$ yr. However, for supermassive BHs we may worry about possible spin growth due to mergers with other BHs and/or accretion. The most likely mechanism to produce fast-spinning BHs is prolonged accretion [33]. Therefore, a conservative assumption to estimate the astrophysical consequences of the instability is to compare the superradiance timescale to the (minimum) timescale over which accretion could spin up the BH. Thin-disk accretion can increase the BH spin from $\tilde{a} = 0$ to $\tilde{a} \simeq 1$ with a corresponding mass increase by a factor $\sqrt{6}$ [34]. If we assume that mass growth occurs via accretion at the Eddington limit, so that the BH mass grows exponentially with e-folding time given by the Salpeter timescale $\tau_{\rm Salpeter} = 4.5 \times 10^7$ yr, then the minimum timescale for the BH spin to grow via thin-disk accretion is comparable to $\tau_{\rm Salpeter}$.

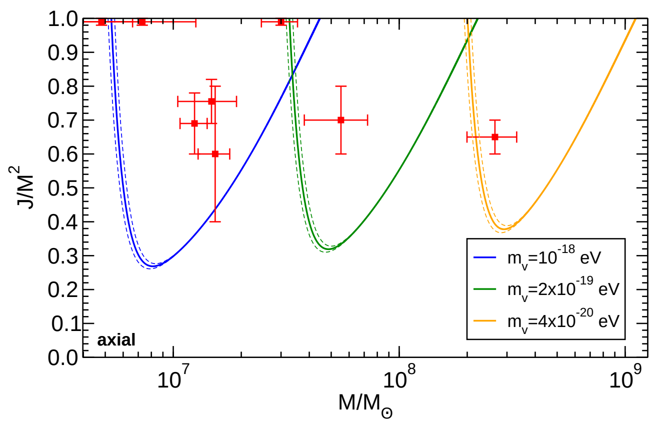

FIG. 1. Contour plots in the BH Regge plane [2] corresponding to an instability timescale shorter than $\tau_{\rm Salpeter}$ for different values of the vector field mass $m_v = \mu\hbar$ and for axial modes with $\ell = m = 1$. Dashed lines bracket our estimated numerical errors, $\gamma_{01} \approx 0.09 \pm 0.03$ in Eq. (7). The experimental points (with error bars) refer to the supermassive BHs listed in Table 2 of [35] and the rightmost point corresponds to the supermassive BH in Fairall 9 [36]. Supermassive BHs lying above each of these curves would be unstable on an observable timescale, and therefore each point rules out a range of Proca field masses.

Brenneman et al. [35] have recently presented a list of eight supermassive BH spin estimates. In order to quantify the dependence of Proca field mass bounds on the mass and spin of supermassive BHs, in Fig. 1 we show exclusion regions in the "BH Regge plane" (cf. Fig. 3 of [2]). To be more specific, we plot contours corresponding to an instability timescale of the order of the Salpeter time for three different masses of the Proca field ($m_v = 10^{-18}$ eV, 2×10^{-19} eV, 4×10^{-20} eV) and for axial modes ($S = 0$). The plot shows that observations of supermassive BHs with $10^6 M_\odot \lesssim M \lesssim 10^9 M_\odot$ and $\tilde{a} \gtrsim 0.3$ would exclude a wide range of vector field masses.

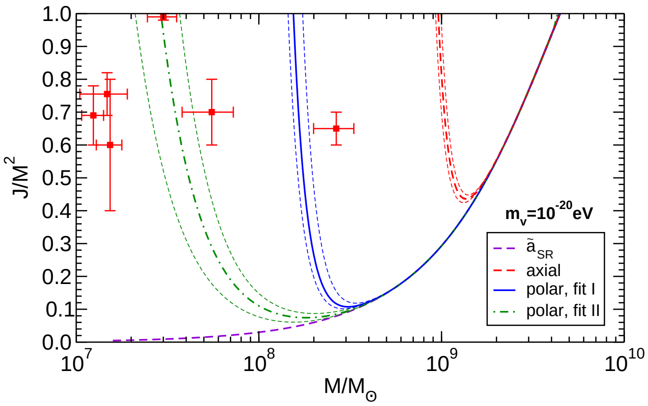

FIG. 2. Comparison between axial and polar $\ell = m = 1$ instability windows for $m_v = 10^{-20}$eV. The right boundary of the instability window does not depend on uncertainties in the fits and it is given by $J/M^2 \equiv \tilde{a} \sim \tilde{a}_{\rm SR} \sim 4M\mu/m + \mathcal{O}(\mu^3)$, corresponding to the superradiance threshold [Eq. (6)] when $\omega_R \sim \mu$. For polar modes we show two different fitting functions. Fit I corresponds to Eq. (7) with $\gamma_{-11} \approx 20 \pm 10$, i.e, $M\omega_I = 20(\tilde{a} - 2r_+\mu)(M\mu)^7$. Fit II is given by $M\omega_I \sim (\tilde{a} - \tilde{a}_{\rm SR})\left[\eta_0(M\mu)^{\kappa_0} + \eta_1\tilde{a}(M\mu)^{\kappa_1}\right]$ with $\eta_0 \approx -6.5 \pm 2$, $\eta_1 \approx 2.1 \pm 1$, $\kappa_0 \approx 6.0 \pm 0.1$, $\kappa_1 \approx 5.0 \pm 0.3$. While fit I is physically more appealing [24], fit II does a better job at reproducing our numerical data in the whole instability region.

In Fig. 2 we compare the axial and polar instability windows for $m_v = 10^{-20}$eV, including our estimated errors. In the polar case we use two different fitting functions, to bracket uncertainties. The polar instability window clearly depends on the chosen fitting function, but some general conclusion can be drawn: (i) axial bounds on m_v are typically less stringent than polar bounds by up to one order of magnitude or more; (ii) different fitting functions translate into bounds on m_v that differ by a factor of a few; (iii) the instability windows in the polar case extend down to $\tilde{a} \sim 0.1$. Thus, essentially *any* supermassive BH spin measurement would exclude a considerable range of m_v; (iv) from Eq. (6) it is straighforward to obtain a conservative *upper* bound for the boson mass that can be excluded by a given observation:

$$m_v \lesssim m_v^{(c)} = \frac{m\tilde{a}\hbar}{4M} \sim 3.34 \times 10^{-19}\, m\tilde{a}\frac{10^8 M_\odot}{M}\,\mathrm{eV}\,. \quad (8)$$

This bound is valid independently of uncertainties in the fitting functions.

Thus, *existing* measurements of supermassive BH spins rule out vector field masses in the whole range 10^{-20} eV $\lesssim m_v \lesssim 10^{-17}$ eV. The best bound comes from Fairall 9 [36], for which the axial instability implies a conservative bound (including measurement and numerical errors) $m_v \lesssim 4 \times 10^{-20}$ eV when we compare the instability timescale to the Salpeter time, and $m_v \lesssim 2 \times 10^{-20}$ eV if we do not consider accretion. This result is of great significance, since it is two orders of magnitude more stringent than the current best bound on the photon mass, $m_\gamma < 10^{-18}$ eV [37]. If the largest known supermassive BHs with $M \simeq 2 \times 10^{10} M_\odot$ [38, 39] were confirmed to have nonzero spin, we could get bounds as low as $m_v \lesssim 10^{-22}$ eV.

Nonlinear effects, other couplings. An important ingredient that was not taken into account in our study is the nonlinear evolution of the instability, that can modify the background geometry. Photon self-interactions are very weak, being suppressed by the mass of the electron. Therefore, it is quite likely that the outcome of the instability will be a slow and gradual drainage of the hole's rotational energy. Another important issue is whether the coupling of accreting matter to massive bosons can quench the instability. In principle massive photons (unlike hidden $U(1)$ fields, for which the interaction with matter is very small) can couple strongly to matter. However it is unlikely that this will significantly affect the superradiant instability discussed here, for two reasons: (i) the unstable modes are large-scale coherent modes whose Compton wavelength is of order the BH size or larger, and accretion disks are typically charge-neutral over these lengthscales, so any possible coupling with ordinary neutral matter is incoherent and most likely inefficient; (ii) accretion disks are often localized in the equatorial plane, and therefore they can affect at most some (but not all) unstable modes. The investigation of the superradiant instability in the presence of matter requires further work, but these arguments suggest that estimates in vacuum should be robust. Spin measurements for slowly accreting BHs (such as the BH at the center of our own galaxy) are presumably the most reliable, being rather insensitive to details of the interaction of vector fields with matter.

Conclusions. The results discussed here show that BHs offer the exciting possibility to constrain particle physics and to set stringent upper bounds on the mass of bosonic fields. Our method can be generalized to other fields and applied to other backgrounds, such as Kerr-Newman BHs or higher-dimensional BHs in TeV-gravity scenarios. Numerical simulations of either linearized or nonlinear perturbations of the Kerr geometry are necessary. In numerical studies of superradiant instabilities a linear perturbation analysis proved to be very useful, e.g. for choosing suitable initial data [22]. Numerical methods may give us a better understanding of the maximum instability rate, of the nonlinear evolution and of the end state of the instability. These are clearly important topics: for example, a more accurate quantitative analysis of the polar sector could improve the bounds on the photon mass by by up to one order of magnitude or more.

Acknowledgments. We wish to thank the Axiverse Project members (especially Hideo Kodama) and Antonino Flachi for valuable discussions, and the referees for their useful suggestions. This work was supported by the DyBHo–256667 ERC Starting Grant, the NRHEP–295189 FP7–PEOPLE–2011–IRSES Grant, and by FCT - Portugal through PTDC projects FIS/098025/2008, FIS/098032/2008, CTE-

ST/098034/2008, CERN/FP/123593/2011. A.I. was supported by JSPS Grant-in-Aid for Scientific Research Fund No. 22540299 and No. 22244030. E.B. was supported by NSF Grant No. PHY-0900735 and NSF CAREER Grant No. PHY-1055103. P.P. acknowledges financial support provided by the European Community through the Intra-European Marie Curie contract aStronGR-2011-298297 and the kind hospitality of the University of Rome "Sapienza" and of the International School for Advanced Studies (SISSA) in Trieste. V.C. thanks the Yukawa Institute for Theoretical Physics at Kyoto University for hospitality during the YITP-T-11-08 workshop on "Recent advances in numerical and analytical methods for black hole dynamics".

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