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Through the looking glass into the dark dimension: Searching for bulk black hole dark matter with microlensing of X-ray pulsars

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ABSTRACT

Primordial black holes (PBHs) hidden in the incredible bulk of the dark dimension could escape constraints from non-observation of their Hawking radiation. Since these five-dimensional (5D) PBHs are bigger, colder, and longer-lived than usual 4D PBHs of the same mass M, they could make all cosmological dark matter if $10^{11} \lesssim M/\mathrm{g} \lesssim 10^{21}$, i.e., extending the 4D allowed region far down the asteroid-mass window. We show that these evasive PBHs could be search for by measuring their X-ray microlensing events from faraway pulsars. We also show that future X-ray microlensing experiments will be able to probe the interesting range ($10^{16.5} \lesssim M/\mathrm{g} \lesssim 10^{17.5}$ g) where an all dark matter interpretation in terms of 4D Schwarzschild PBHs is excluded by the non-observation of their Hawking radiation.

1. Introduction

The dark dimension, an innovative five-dimensional (5D) scenario that has a compact space with characteristic length-scale in the micron range, provides a stepping stone to address the cosmological hierarchy problem [1]. This is because the anti-de Sitter distance conjecture in de Sitter space [2] connects the size of the compact space $R_\perp$ to the dark energy scale $\Lambda^{1/4}$ via $R_\perp \sim \lambda \Lambda^{1/4}$, where $\Lambda \sim 10^{-122} M_p^4$ is the cosmological constant, M_p the Planck mass, and the proportionality factor is estimated to be within the range $10^{-1} < \lambda < 10^{-4}$ [1].

Primordial black holes (PBHs), presumably formed through the collapse of sufficiently sizable overdensities originated by an enhancement in the comoving curvature perturbation power spectrum at small scales during inflation [3–5], may be one of the most interesting candidates of dark matter [6–8]. These intriguing dark matter candidates are particular important in the dark dimension scenario, because 5D PBH are bigger, colder, and longer-lived than usual 4D PBHs of the same mass M [9–11].[1] Furthermore, we have recently shown that a back reaction of Hawking evaporation process is to kick 5D Schwarzschild PBHs out of the brane [18]. As a consequence, 5D PBHs could make all cosmological dark matter if $10^{11} \lesssim M/\mathrm{g} \lesssim 10^{21}$, i.e., well down into the asteroid-mass window. In this paper we investigate a potential signal that may allow us to unmask these as yet evasive 5D PBHs.

An astonishing coincidence is that $R_\perp$ corresponds approximately to the wavelength of visible light. This implies that the Schwarzschild radius of black holes perceiving the dark dimension is well below the wavelength of light. For point-like lenses, this is precisely the critical length where geometric optics breaks down and the effects of wave optics suppress the magnification [19–21], obstructing the sensitivity to 5D PBH microlensing signals. Nevertheless, it was pointed out in [22] that X-ray pulsars, with photon energies above 1 keV, are good candidate sources to search for microlensing of PBHs with $M \lesssim 10^{21}$ g. In this paper we reexamine this idea focussing attention on 5D PBHs perceiving the dark dimension.

The layout is as follows. In Section 2 we review the generalities of gravitational microlensing. In Section 3 we focus attention on the particulars of microlensing of X-ray pulsars, which is relevant to unmask PBHs perceiving the dark dimension. The paper wraps up in Section 4 with some conclusions.

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[1] The dark dimension scenario provides a specific realization of the dynamical dark matter framework [12] with other interesting dark matter candidates which have been discussed in [13–17].

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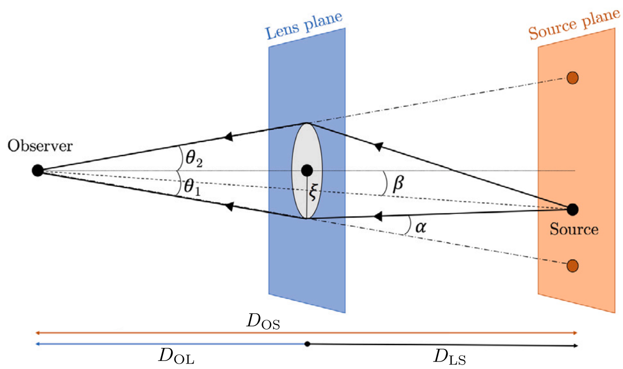

Fig. 1. The geometry of the microlensing setup for a point-like lens, seen to produce two images after deflection [25].

2. Point-lens microlensing basics

It has long been known that gravitational lensing becomes observable when a massive object is located between a light-emitting source and an observer [23]. This is because the gravitational potential of the central object acts as a lens that bends the path of the light rays coming from the source, by bending the spacetime through which the photon travels. If the central object is a Schwarzschild black hole, described by the line element

$$ds^2 = -f(r)\,dt^2 + f^{-1}(r)\,dr^2 + r^2\,d\Omega_2^2\,, \tag{1}$$

the deflection angle of a light ray is estimated to be

$$\alpha = \frac{4GM}{r_0} + 4\left(\frac{GM}{r_0}\right)^2\left(\frac{15\pi}{16} - 1\right) + \mathcal{O}\left(\epsilon^3\right)\,, \tag{2}$$

where $f(r) = 1 - 2GM/r$ is the blackening function, M is the mass of the black hole (i.e. the lens), $G = M_p^{-2}$ is Newton's gravitational constant, $d\Omega_2^2$ is the metric of a 2-dimensional unit sphere, and r_0 is the closest distance of approach to the lens in the lens-Earth-source plane, with $\epsilon = GM/r_0$, and where all measurements are taken in the observer's reference frame [24]. To develop some sense for the orders of magnitude involved, we consider the set up shown in Fig. 1 for a point-like lens. In the absence of the lens, we would then observe the source at an angular position β on the source plane, but because of the deflection, we actually observe it at θ. Since the observer distances to the lens $D_{\rm OL}$ and source $D_{\rm OS}$ are very large in relation to the deflection angle, the distances of the source and image from the optical axis in the source plane are estimated using the small angle approximation

$$\theta D_{\rm OS} = \beta D_{\rm OS} + \alpha D_{\rm LS}\,, \tag{3}$$

where $D_{\rm LS}$ is the distance from the lens to the source [19].

The impact parameter of the unperturbed light ray is found to be $\xi = r_0(1 - 2GM/r_0)^{-1/2}$. In the point-lens approximation, however, $\xi \equiv \theta D_{\rm OL} \sim r_0$, and so (3) can be rewritten as

$$\theta^2 = \beta\theta + \theta_E^2\,, \tag{4}$$

where

$$\theta_E = \sqrt{4GM\,\frac{D_{\rm LS}}{D_{\rm OL}\,D_{\rm OS}}} \tag{5}$$

is the angular size of the Einstein ring which is formed when the source is perfectly aligned behind the lens, i.e., when $\beta = 0$, and where we have neglected terms $\mathcal{O}(\epsilon^2)$. θ_E in turn defines the Einstein radius $r_E = D_{\rm OL}\theta_E$ on the lens plane. Note that for $\beta \sim 0$, the impact parameter $\xi \sim r_E$ is the characteristic scale of the source-lens-observer system. All in all, the isolated Schwarzschild black hole will split a point source into two images with angular positions defined implicitly by

$$\theta_{1,2} = \frac{1}{2}\left(\beta \pm \sqrt{\beta^2 + 4\theta_E^2}\right)\,. \tag{6}$$

The two images are on opposite sides of the source, with one image inside the Einstein ring and the other outside. As the source moves away from the lens (i.e. as β increases), one of the images approaches the lens and becomes very faint, whereas the other image approaches the true position of the source and asymptotes to its unlensed flux.

A point worth noting at this juncture is that if a source is closer than θ_E in separation from a lensing PBH on the sky, the source is multiply imaged by its lensing. Nevertheless, the separation between multiple images is too small to be resolved by optical telescopes. What is observed instead in the so-called "microlensing phenomenon" is the temporary magnification of the total flux of two images relative to that of the unlensed source. Now, if a gravitational field were to deflect a light ray this would produce a change in the cross-section of a bundle of rays. However, according to Liouville's theorem the phase space density of photons must be conserved, which implies that gravitational lensing should preserve the surface brightness of the source and should only change its apparent surface area. In other words, the magnification of an image becomes the ratio of the solid angles of the image and of the unlensed source (at the observer position). If the central object is spherically symmetric, the magnification factor is found to be

$$\mu = \frac{\sin\theta\;d\theta}{\sin\beta\;d\beta} \simeq \frac{\theta}{\beta}\;\frac{d\theta}{d\beta}\,, \tag{7}$$

where the sign of the magnification gives the parity of the particular image. Substituting β from the point-lens Eq. (4) into (7), it follows the magnifications of the two images,

$$\mu_{1,2} = \left[1 - \left(\frac{\theta_E}{\theta_{1,2}}\right)^4\right]^{-1} = \frac{u^2+2}{2u\sqrt{u^2+4}} \pm \frac{1}{2}\,, \tag{8}$$

where u is the angular separation of the source from the point mass in units of the Einstein angle, $u = \beta/\theta_E$. Note that $\theta_2 < \theta_E$ implies $\mu_2 < 0$, and so the magnification of the image which is inside the Einstein ring is negative implying that this image has its parity flipped with respect to the source. The net magnification of flux in the two images is obtained by adding the absolute magnifications,

$$\mu_{\rm tot}(u) = \mu_1 + \mu_2 = \frac{u^2+2}{u\sqrt{u^2+4}}\,. \tag{9}$$

Note that when the source lies on the Einstein radius, we have $\beta = \theta_E$ and u = 1, so that the total magnification becomes $\mu_{\rm tot} = 1.17 + 0.17 = 1.34$.

Now, the source and the lens have a relative motion with respect to the observer. As a consequence, the observed flux of a source varies with time, yielding a characteristic light curve of the observed source flux. Such a unique light curve allows a microlensing event to be identified from the observation of other variable sources. A typical timescale of the microlensing light curve can be estimated from a crossing time of the Einstein radius for a lensing PBH with respect to a distant source,

$$t_E = r_E/v\,, \tag{10}$$

where v is the relative velocity for a observer-lens-source system [26].

Microlensing surveys are typically sensitive to stars that are $10 \lesssim D_{\rm OS}/{\rm kpc} \lesssim 1000$ away and to transit times of minutes to years. For example, the Subaru-HSC instrument was sensitive to the short transit times of light PBHs reaching sensitivities to constrain an all dark matter PBH interpretation for masses $M \gtrsim 10^{22}$ g [19–21]. However the Subaru-HSC survey was limited by two effects as the PBH mass is reduced:

- the finite apparent size of the source stars in its sky target, M31, being larger than the apparent Einstein radius, yielding poorly focused light;

- transition from geometric to wave optics as the PBH Schwarzschild radius becomes comparable to or smaller than the wavelength of optical light.

Indeed, the wave effects characterized by

$$w = \frac{8\pi G M}{\lambda} = 0.3 \left(\frac{M}{10^{22}\ \mathrm{g}}\right)\left(\frac{\lambda_0}{621\ \mathrm{nm}}\right)^{-1}, \tag{11}$$

become important when $w \lesssim 1$, where λ_0 is the characteristic wavelength of light in an observation [27]; the default choice in the normalization of (11) corresponds to a central wavelength of r-band in the Subaru telescope. As it was first pointed out in [22], consideration of sources emitting in the X-ray spectrum could allow us to search deep into the PBH low-mass window. It is this that we now turn to study.

3. Microlensing of black holes perceiving the dark dimension

It is a known fact that M_p is the natural cutoff of quantum gravity. However, if there were light species of particles in the theory, then the consistency of black hole entropy with the effective field theory description would demand a breakdown of the classical picture at a lower scale,

$$M_* = m_{\rm KK}^{(d-4)/(d-2)} M_p^{2/(d-2)}, \tag{12}$$

dubbed the species scale, where d is the spacetime dimension and $m_{\rm KK} \sim R_\perp^{-1}$ [28,29]. Now, despite the fact that M_* is motivated by the emergence of the tower of light states, curiously the mass scale of this tower, $m_{\rm KK}$, does not seem to be directly captured by M_* [30].

One way the lower-dimensional theory can find out about the $m_{\rm KK}$ scale is through the study of black holes [10,11,30]. When black holes get smaller than $R_\perp$, the black hole becomes thermodynamically unstable because of the Gregory–Laflamme transition [31]. Black holes with $r_s \gg R_\perp$ are actually black branes wrapped around the extra dimensions, but the ones with $r_s \ll R_\perp$ are localized in the extra dimensions. This transition to a new black hole solution marks a new scale $\Lambda_{\rm BH}$ in the lower-dimensional theory [30]. In this section we show that Earth-based microlensing experiments are insensitive to the $\Lambda_{\rm BH}$ scale, independently of the black hole mass M.

Throughout we assume that the Standard Model fields are confined to a D-brane and only gravity spills into the compact space of dimension $(d-4)$ [32].

3.1. Microlensing of bulk black holes

The Gregory–Laflamme transition induces a change in the scaling of the black hole's Schwarzschild radius r_s and its Hawking temperature T_H [10]. If the black hole is spherically symmetric and $r_s \ll R_\perp$, then it can be treated as a flat d-dimensional object with line element given by

$$ds^2 = -U(r)\ dt^2 + U^{-1}(r)\ dr^2 + r^2\ d\Omega^2_{d-2}, \tag{13}$$

where $U(r) = 1 - (r_s/r)^{d-3}$ is the blackening function,

$$d\Omega^2_{d-2} = d\chi_2^2 + \prod_{i=2}^{d-2} \sin^2\chi_i\ d\chi^2_{i+1} \tag{14}$$

is the metric of a $(d-2)$-dimensional unit sphere, and

$$r_s = \frac{1}{M_*}\left[\frac{M}{M_*}\right]^{1/(d-3)} \left[\frac{8\,\Gamma\left(\frac{d-1}{2}\right)}{(d-2)\ \pi^{(d-3)/2}}\right]^{1/(d-3)}, \tag{15}$$

and where $\Gamma(x)$ is the Gamma function [33–35]. The d-dimensional black hole behaves like a thermodynamic system [36], with temperature $T_H \sim (d-3)/(4\pi r_s)$ and entropy $S = (4\pi M r_s)/(d-2)$ [37]. If the black hole is localized on the brane, Hawking evaporation [38,39] proceeds dominantly through emission of Standard Model fields [40]. However, we have recently shown that the recoil effect due to graviton emission imparts the black hole a relative kick velocity with respect to the brane, allowing Schwarzschild black holes to escape into the bulk [18]. Because the escape from the brane is almost instantaneous, 5D Schwarzschild black holes evaporate almost entirely into gravitons in the bulk, and so can evade constraints from the non-observation of Hawking radiation in: *(i)* the extragalactic γ-ray background [41], *(ii)* the cosmic microwave background [42], *(iii)* the 511 keV γ-ray line [43–46], *(iv)* the EDGES 21-cm signal [47,48], *(v)* Lyman-α forest [49], and *(vi)* the MeV Galactic diffuse emission [50–52].

The deflection angle of a light ray expected to be induced by a d-dimensional Schwarzschild black hole localized on the brane is estimated to be

$$\alpha_d = \frac{2(d-2)M}{M_*^{d-2}\ r_0^{d-3}}\ {}_2F_1\left(\frac{1}{2},\kappa;\frac{3}{2};1\right) + \mathcal{O}\left(\frac{M^2}{M_*^d r_0^{d-2}}\right), \tag{16}$$

where ${}_2F_1(a,b,;c;z)$ is the Gaussian hypergeometric function, with $\kappa = 1/2-(d-3)/2$ [53]. A straightforward calculation shows that for $d=4$, (16) reduces to the first term in (2). For $d=5$,

$$\alpha_5 = \frac{3}{2}\ \frac{\pi}{M_*^3}\ \frac{M}{r_0^2} + \mathcal{O}\left(\frac{M^2}{M_*^5 r_0^3}\right), \tag{17}$$

which implies that $\alpha_5 < \alpha_4$ for fixed M and r_0, with $r_s < r_0 < R_\perp$. This was interpreted in [53] as the need for advanced sensitive detection devices to observe lensed images influenced by the dark dimension.

The previous statement should be revised with caution. It is clear that while the source is on the brane, in the dark dimension scenario the lens has an extension into the bulk. Therefore, the light emitted by the source could be a micron away from the lens 4D position. The scale that counts in addressing whether the 5D geometry of the lens would bring important corrections to the source magnification $\mu_{\rm tot}$ is actually the size ξ of the lens involved in the deflection. As we have stressed in Section 2, in the point-lens approximation of Fig. 1, the characteristic scale of the source-lens-observer system is $\xi \gtrsim r_E$.

At this point a reality check is in order. A source could be a good object to undergo microlensing if it is a long distance away from the telescopes (viz., the Earth), because this would increase the optical depth and/or the number of possible lensing events. Following [22], we consider the X-ray pulsar SMC X-1 in the Small Magellanic Cloud at a distance $D_{\rm OS} = 65$ kpc [54]. For such a distance, the Einstein radius of a $M \sim 10^{17}$ g black hole is given by

$$\begin{aligned} r_E &= \sqrt{4GMx(1-x)D_{\rm OS}} \\ &= 12\ \mathrm{km}\left[\left(\frac{M}{10^{17}\ \mathrm{g}}\right)\left(\frac{D_{\rm OS}}{65\ \mathrm{kpc}}\right)\left(\frac{x(1-x)}{1/4}\right)\right]^{1/2} \end{aligned} \tag{18}$$

where $x = D_{\rm OL}/D_{\rm OS}$. Since $r_E \gg R_\perp$ is the closest distance that source photons get to the lens when it lies directly along the line of sight, we conclude that corrections to $\mu_{\rm tot}$ due to the 5D geometry of the lens can be safely neglected. In other words, for $r \gg R_\perp$ the gravitational potential of a 5D black hole falls as $1/r$ and therefore is indistinguishable from that of a 4D black hole of the same mass. This implies that for Earth-based experiments the microlensing signal of a 5D black hole would be indistinguishable from that of a 4D black hole of the same mass, and this is independent on whether the black hole is localized on the brane or propagates through the bulk.

Since line-of-sight distances of interest are usually much larger than transverse distances, X-ray microlensing events can be pictured as projections on the lens-containing transverse plane, *viz.* the lens plane. The source radius in the lens plane is estimated to be

$$\begin{aligned} a_s(x) &= \frac{xR_s}{r_E(x)} \sim 0.8 \times \left(\frac{x}{\sqrt{x(1-x)}}\right)\left(\frac{R_S}{20\ \mathrm{km}}\right) \\ &\quad\times \left(\frac{D_{\rm OS}}{65\ \mathrm{kpc}}\right)^{-1}\left(\frac{M}{10^{17}\ \mathrm{g}}\right)^{-1/2}, \end{aligned} \tag{19}$$

where following [22] we have taken a fiducial source size $10 \lesssim R_S/\mathrm{km} \lesssim 20$. From (19) we infer that X-ray pulsars could overcome finite source size effects for $M \sim 10^{17}$ g if $x \lesssim 0.6$.

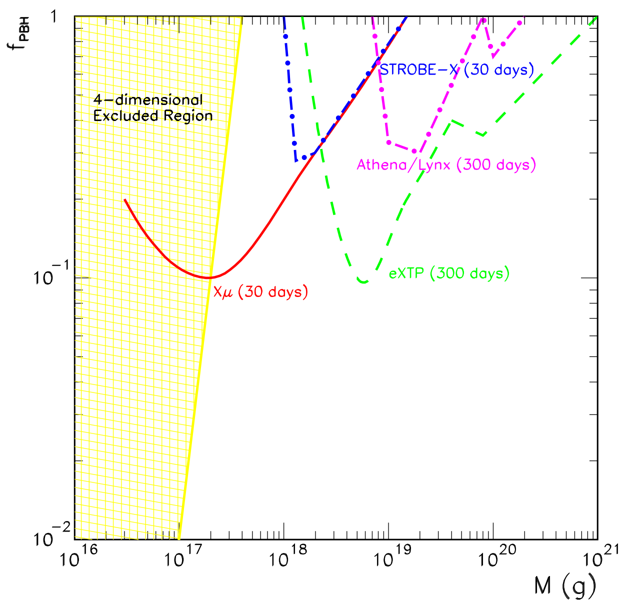

Fig. 2. Compilation of projected sensitivities on the PBH dark matter fraction $f_{\rm PBH}$ as a function of M for future X-ray telescopes, assuming a monochromatic mass function and including wave optics effects. The sensitivity curves have been taken from [22,56]. For comparison, the shaded region shows constraints on the 4D $f_{\rm PBH}$ from the non-observation of Hawking radiation in: *(i)* the extragalactic γ-ray background [41], *(ii)* the cosmic microwave background [42], *(iii)* the 511 keV γ-ray line [43–46], *(iv)* the EDGES 21-cm signal [47,48], *(v)* Lyman-α forest [49], and *(vi)* the MeV Galactic diffuse emission [50–52].

Now, wave effects becomes important when $w \lesssim 1$ [27] or equivalently for photon energies

$$E_\gamma \lesssim \frac{1}{4GM} = 660 \text{ keV} \left(\frac{M}{10^{17}\text{ g}}\right)^{-1} . \tag{20}$$

Thus, for a source emitting primarily with X-ray energy of $1 \lesssim E_\gamma/\text{keV} \lesssim 10$ [54], wave effects must be taken into account in order to probe a lower PBH mass of about 10^{17} g.

The magnification, including the wave optics effect, is given by

$$\mu_{\rm tot}(w,u) = \frac{\pi w}{1-e^{-\pi w}} \left| {}_1F_1\left(\frac{i}{2}w, 1; \frac{i}{2}wu^2\right)\right|^2 \tag{21}$$

where ${}_1F_1(a;b;z)$ is the confluent hypergeometric function [55]. For $u = 0$ the magnification has a maximum

$$\mu_{\rm tot}(w,u)\Big|_{\rm max} = \mu_w(w,0) = \frac{\pi w}{1-e^{-\pi w}} . \tag{22}$$

It is easily seen from (22) that for microlensing events with $w \ll 1$, the maximum magnification

$$\mu_{\rm tot}(w,u)\Big|_{\rm max} = 1 + \frac{\pi w}{2} , \tag{23}$$

would be significantly reduced as compared to the maximum magnification $\mu_{\rm tot}(u) \to \infty$ in the geometric optics approximation.[2] This is because the gravitational potential induced by extremely light PBHs is too weak to bent the path of the photons.

The mass distribution of PBHs is usually characterized by the mass function

$$\psi(M) = \frac{M}{\rho_{\rm CDM}} \frac{dn_{\rm PBH}}{dM} , \tag{24}$$

where $dn_{\rm PBH}$ is the number density of PBHs within the mass range $(M, M + dM)$, and $\rho_{\rm CDM}$ is the energy density of cold dark matter [57].

Integrating $\psi(M)$ gives the total fraction of dark matter in PBHs,

$$f_{\rm PBH} \equiv \frac{\rho_{\rm PBH}}{\rho_{\rm CDM}} = \int \psi(M)\, dM , \tag{25}$$

where $\rho_{\rm PBH} = \int M\, dn_{\rm PBH}$ is the energy density of PBHs. If all of the dark matter were made of PBH, we would have $f_{\rm PBH} = 1$. In Fig. 2 we show a compilation of the projected sensitivity to $f_{\rm PBH}$ of future X-ray telescopes, including Athena [58], Lynx [59], and eXTP for a 300-day observation [60], as well as STROBE-X [61], and $X\mu$ for a 30-day observation [56]. We can see that these instruments will break down into the mass range relevant for bulk black hole dark matter, including the interesting region in which a 4D PBH all dark matter interpretation has been excluded due to the non-observation of Hawking radiation.

3.2. Microlensing of near-extremal black holes

According to the no-hair theorem [62], 4D black hole geometries in asymptotically flat spacetimes eventually settled down to Kerr–Newman solutions [63,64], which are characterized by three measurable parameters: the mass M, the angular momentum $\vec{J}$ and the electric charge Q. An important feature of Kerr–Newman black holes is that the three parameters are not all independent from each other. Actually, for a given set of parameters there is a minimal extremal mass M_e satisfying

$$M^2 \geq M_e^2 = \left(\frac{M_p^2 J}{M}\right)^2 + (M_p Q)^2 , \tag{26}$$

where Q is measured in units in which the Coulomb force between two charges separated by a distance d has magnitude $F = Q^2/d^2$ and $J = |\vec{J}|$ [65]. If the parameters saturate this bound the black hole is dubbed extremal, and if the parameters are close to saturating it near-extremal. Before proceeding we pause to note that:

- When the black hole mass is tuned below $M_{\rm ext}$, the event horizon disappears leaving behind a naked singularity, which violates the cosmic censorship conjecture [66].
- When a black hole reaches its extreme limit, the thermal description breaks down, and it cannot continue to evaporate by emitting (uncharged) elementary particles.
- When black holes are near-extremal their evaporation temperature decreases, and consequently so does their luminosity. Thus, near-extremal black holes can also evade constraints from the non-observation of Hawking radiation [67,68].

It has long been suspected that any electromagnetic charge or spin would be lost very quickly by any 4D black hole population of primordial origin. On the one hand, the electromagnetic charge of a black hole is spoiled by the Schwinger effect [69], which allows pair-production of electron-positron pairs in the strong electric field outside the black hole, leading to the discharge of the black hole and subsequent evaporation [70,71]. On the other hand, a rapidly rotating black hole spins down to a nearly non-rotating state before most of its mass has been given up, and therefore it does not approach to extremal when it evaporates [72]. All in all, near-extremal primordial Kerr–Newman black holes are not expected to prevail in the universe we live in.

An alternative interesting possibility is to envision a scenario where the black hole is charged under a generic unbroken $U(1)$ symmetry (dark photon), whose carriers (dark electrons with a mass m_e' and a gauge coupling e') are always much heavier than the temperature of the black hole [73]. This implies that the charge Q does not get evaporated away from the black hole and remains therefore constant. Strictly speaking, the pair production rate per unit volume from the Schwinger effect can be slowed down by arbitrarily decreasing e', whereas the weak gravity conjecture (WGC) imposes a constraint on the charge per unit mass; namely, for each conserved gauge charge there must be a sufficiently light charge carrier such that

$$e'q/m_{e'} \geq \sqrt{4\pi}\sqrt{(d-3)/(d-2)}\; \bar{M}_p^{-(d-2)/2} , \tag{27}$$

[2] Note that $\mu_{\rm tot}(u,w) = 1$ means no lensing magnification.

where q is the integer-quantized electric charge of the particle and $\bar{M}_p = M_p/(8\pi)$ is the reduced Planck mass [74,75]. Setting $e = e' = \sqrt{4\pi\alpha}$ the (4D) Schwinger effect together with the WGC lead to a bound on the minimum black hole mass of near extremal black holes with evaporation time longer than the age of Universe, $M_{\rm ne} \gtrsim 5 \times 10^{15}$ g$(m_{e'}/10^9~{\rm GeV})^{-2}$ [73].

Since we have seen that microlensing experiments cannot distinguish 4D from 5D black holes, in what follows we consider the line element (1) with blackening function

$$f(r) = 1 - \frac{2M}{M_p^2\, r} + \frac{Q^2}{M_p^2\, r^2}\,. \tag{28}$$

The deflection angle of a light ray expected to be induced by the geometry given in (28) has been computed in [76] and is given by

$$\alpha_Q = \frac{4M}{M_p^2\, r_0} + \left(\frac{15}{16}\pi - 1\right)\frac{4M^2}{M_p^4 r_0^2} - \frac{3}{4}\pi\frac{Q^2}{M_p^2 r_0^2} + \cdots\,. \tag{29}$$

We can see that the effect of Q is to slightly reduce the second order correction to the deflection angle. Therefore, we can conclude that the sensitivity of future X-ray microforlensing experiments to near-extremal black holes is actually shown in Fig. 2.

In summary, evaporation constraints on $f_{\rm PBH}$ can be substantially altered when moving away from the Schwarzschild picture. (Other PBH models where Hawking radiation can be slowed down have been recently discussed in [77,78].) Microlensing experiments of X-ray pulsars are well positioned to uncover these models.

4. Conclusions

We have investigated the potential of searching for bulk black hole dark matter through microlensing events using future X-ray telescopes. We have shown that for these telescopes the microlensing signal of a 5D black hole would be indistinguishable from that of a 4D black hole of the same mass, and this is independent on whether the black hole is localized on the brane or propagates through the bulk. We have demonstrated that future instruments observing microlensing events from X-ray pulsars will probe the mass range relevant for bulk black hole dark matter, including the interesting region in which a 4D PBH all dark matter interpretation has been excluded due to the non-observation of Hawking radiation.

We end with an observation. PBHs may experience a memory burden effect, which splits the evaporation process into two distinct phases: semiclassical and quantum [79,80]. If this were the case, then the minimum black hole mass allowing a PBH all-dark-matter interpretation would also be relaxed [81–84]. The quantum decay rate has an additional suppression factor compared to the Hawking decay rate, which in the most realistic scenario scales with the inverse of the black hole entropy. For $d = 4$ and a quantum decay rate $\Gamma \sim T_H/S$, the memory burden effect opens a new window in the mass range $10^9 \lesssim M/{\rm g} \lesssim 10^{10}$ [83]. In contrast to what is stated in [56], using (19) we argue that due to the extended nature of the pulsar emission region X-ray microlensing experiments would be inappropriate to search for signals of the memory burden effect.

CRediT authorship contribution statement

Luis A. Anchordoqui: Formal analysis, Conceptualization, Writing – original draft. **Ignatios Antoniadis:** Formal analysis, Conceptualization, Writing – original draft. **Dieter Lüst:** Formal analysis, Conceptualization, Writing – original draft. **Karem Peñaló Castillo:** Formal analysis, Conceptualization, Writing – original draft.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

Data availability

No data was used for the research described in the article.

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References

[1] M. Montero, C. Vafa, I. Valenzuela, The dark dimension and the swampland, JHEP 02 (022) (2023) http://dx.doi.org/10.1007/JHEP02(2023)022, arXiv:2205.12293, [hep-th].
[2] D. Lüst, E. Palti, C. Vafa, Ads and the swampland, Phys. Lett. B 797 (2019) 134867, http://dx.doi.org/10.1016/j.physletb.2019.134867, arXiv:1906.05225, [hep-th].
[3] Y.B. Zel'dovich, I.D. Novikov, The hypothesis of cores retarded during expansion and the hot cosmological model, Soviet Astron. AJ (Engl. Transl.) 10 (1967) 602.
[4] S. Hawking, Gravitationally collapsed objects of very low mass, Mon. Not. Roy. Astron. Soc. 152 (1971) 75.
[5] B.J. Carr, S.W. Hawking, Black holes in the early universe, Mon. Not. Roy. Astron. Soc. 168 (1974) 399–415.
[6] B. Carr, F. Kuhnel, Primordial black holes as dark matter: recent developments, Ann. Rev. Nucl. Part. Sci. 70 (2020) 355–394, http://dx.doi.org/10.1146/annurev-nucl-050520-125911, arXiv:2006.02838, [astro-ph.CO].
[7] A.M. Green, B.J. Kavanagh, Primordial black holes as a dark matter candidate, J. Phys. G 48 (4) (2021) 043001, http://dx.doi.org/10.1088/1361-6471/abc534, arXiv:2007.10722, [astro-ph.CO].
[8] P. Villanueva-Domingo, O. Mena, S. Palomares-Ruiz, A brief review on primordial black holes as dark matter, Front. Astron. Space Sci. 8 (2021) 87, http://dx.doi.org/10.3389/fspas.2021.681084, arXiv:2103.12087, [astro-ph.CO].
[9] L.A. Anchordoqui, I. Antoniadis, D. Lüst, Dark dimension, the swampland, and the dark matter fraction composed of primordial black holes, Phys. Rev. D 106 (8) (2022) 086001, http://dx.doi.org/10.1103/PhysRevD.106.086001, arXiv:2206.07071, [hep-th].
[10] L.A. Anchordoqui, I. Antoniadis, D. Lüst, The dark universe: primordial black hole ⇋ dark graviton gas connection, Phys. Lett. B 840 (2023) 137844, http://dx.doi.org/10.1016/j.physletb.2023.137844, arXiv:2210.02475, [hep-th].
[11] L.A. Anchordoqui, I. Antoniadis, D. Lüst, More on black holes perceiving the dark dimension, Phys. Rev. D 110 (1) (2024) 015004, http://dx.doi.org/10.1103/PhysRevD.110.015004, arXiv:2403.19604, [hep-th].
[12] K.R. Dienes, B. Thomas, Dynamical dark matter i: theoretical overview, Phys. Rev. D 85 (2012) 083523, http://dx.doi.org/10.1103/PhysRevD.85.083523, arXiv:1106.4546, [hep-ph].
[13] E. Gonzalo, M. Montero, G. Obied, C. Vafa, Dark dimension gravitons as dark matter, JHEP 11 (2023) 109, http://dx.doi.org/10.1007/JHEP11(2023)109, arXiv:2209.09249, [hep-ph].
[14] L.A. Anchordoqui, I. Antoniadis, D. Lüst, Fuzzy dark matter and the dark dimension, Eur. Phys. J. C 84 (3) (2024) 273, http://dx.doi.org/10.1140/epjc/s10052-024-12622-y, arXiv:2307.01100, [hep-ph].
[15] J.A.P. Law-Smith, G. Obied, A. Prabhu, C. Vafa, Astrophysical constraints on decaying dark gravitons, JHEP 06 (2024) 047, http://dx.doi.org/10.1007/JHEP06(2024)047, arXiv:2307.11048, [hep-ph].
[16] G. Obied, C. Dvorkin, E. Gonzalo, C. Vafa, Dark dimension and decaying dark matter gravitons, Phys. Rev. D 109 (6) (2024) 063540, http://dx.doi.org/10.1103/PhysRevD.109.063540, arXiv:2311.05318, [astro-ph.CO].
[17] N. Gendler, C. Vafa, Axions in the dark dimension, arXiv:2404.15414, [hep-th].
[18] L.A. Anchordoqui, I. Antoniadis, D. Lüst, K.P. Castillo, Bulk Black Hole Dark Matter, arXiv:2407.21031, [hep-th].
[19] N. Matsunaga, K. Yamamoto, The finite source size effect and the wave optics in gravitational lensing, JCAP 01 (2006) 023, http://dx.doi.org/10.1088/1475-7516/2006/01/023, arXiv:astro-ph/0601701, [astro-ph].
[20] S. Sugiyama, T. Kurita, M. Takada, On the wave optics effect on primordial black hole constraints from optical microlensing search, Mon. Not. Roy. Astron. Soc. 493 (3) (2020) 3632–3641, http://dx.doi.org/10.1093/mnras/staa407, arXiv:1905.06066, [astro-ph.CO].
[21] D. Croon, D. McKeen, N. Raj, Z. Wang, Subaru-hsc through a different lens: microlensing by extended dark matter structures, Phys. Rev. D 102 (8) (2020) 083021, http://dx.doi.org/10.1103/PhysRevD.102.083021, arXiv:2007.12697, [astro-ph.CO].

[22] Y. Bai, N. Orlofsky, Microlensing of X-ray pulsars: a method to detect primordial black hole dark matter, Phys. Rev. D 99 (12) (2019) 123019, http://dx.doi.org/10.1103/PhysRevD.99.123019, arXiv:1812.01427, [astro-ph.HE].
[23] A. Einstein, Lens-like action of a star by the deviation of light in the gravitational field, Science 84 (1936) 506–507, http://dx.doi.org/10.1126/science.84.2188.506.
[24] K.S. Virbhadra, G.F.R. Ellis, Schwarzschild black hole lensing, Phys. Rev. D 62 (2000) 084003, http://dx.doi.org/10.1103/PhysRevD.62.084003, arXiv:astro-ph/9904193, [astro-ph].
[25] D. Croon, D. McKeen, N. Raj, Gravitational microlensing by dark matter in extended structures, Phys. Rev. D 101 (8) (2020) 083013, http://dx.doi.org/10.1103/PhysRevD.101.083013, arXiv:2002.08962, [astro-ph.CO].
[26] H. Niikura others, Microlensing constraints on primordial black holes with subaru/hsc andromeda observations, Nature Astron. 3 (6) (2019) 524–534, http://dx.doi.org/10.1038/s41550-019-0723-1, arXiv:1701.02151, [astro-ph.CO].
[27] A. Gould, Femtolensing of gamma-ray bursters, Astrophys. J. Lett. 386 (1992) L5, http://dx.doi.org/10.1086/186279.
[28] G. Dvali, Black holes and large N species solution to the hierarchy problem, Fortsch. Phys. 58 (2010) 528–536, http://dx.doi.org/10.1002/prop.201000009, arXiv:0706.2050, [hep-th].
[29] G. Dvali, M. Redi, Black hole bound on the number of species and quantum gravity at lhc, Phys. Rev. D 77 (2008) 045027, http://dx.doi.org/10.1103/PhysRevD.77.045027, arXiv:0710.4344, [hep-th].
[30] A. Bedroya, C. Vafa, D.H. Wu, The tale of three scales: the Planck, the species, and the black hole scales, arXiv:2403.18005, [hep-th].
[31] R. Gregory, R. Laflamme, Black strings and p-branes are unstable, Phys. Rev. Lett. 70 (1993) 2837–2840, http://dx.doi.org/10.1103/PhysRevLett.70.2837, arXiv:hep-th/9301052, [hep-th].
[32] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, G.R. Dvali, New dimensions at a millimeter to a fermi and superstrings at a tev, Phys. Lett. B 436 (1998) 257–263, http://dx.doi.org/10.1016/S0370-2693(98)00860-0, arXiv:hep-ph/9804398, [hep-ph].
[33] F.R. Tangherlini, Schwarzschild field in n dimensions and the dimensionality of space problem, Nuovo Cim. 27 (1963) 636–651, http://dx.doi.org/10.1007/BF02784569.
[34] R.C. Myers, M.J. Perry, Black holes in higher dimensional space–times, Annals Phys. 172 (1986) 304, http://dx.doi.org/10.1016/0003-4916(86)90186-7.
[35] P.C. Argyres, S. Dimopoulos, J. March-Russell, Black holes and submillimeter dimensions, Phys. Lett. B 441 (1998) 96–104, http://dx.doi.org/10.1016/S0370-2693(98)01184-8, arXiv:hep-th/9808138, [hep-th].
[36] S.W. Hawking, Black holes and thermodynamics, Phys. Rev. D 13 (1976) 191–197, http://dx.doi.org/10.1103/PhysRevD.13.191.
[37] L.A. Anchordoqui, J.L. Feng, H. Goldberg, A.D. Shapere, Black holes from cosmic rays: probes of extra dimensions and new limits on tev scale gravity, Phys. Rev. D 65 (2002) 124027, http://dx.doi.org/10.1103/PhysRevD.65.124027, arXiv:hep-ph/0112247, [hep-ph].
[38] S.W. Hawking, Black hole explosions, Nature 248 (1974) 30–31, http://dx.doi.org/10.1038/248030a0.
[39] S.W. Hawking, Particle creation by black holes, Commun. Math. Phys. 43 (1975) 199–220, http://dx.doi.org/10.1007/BF02345020, [erratum: Commun. Math. Phys. 46 (1976) 206].
[40] R. Emparan, G.T. Horowitz, R.C. Myers, Black holes radiate mainly on the brane, Phys. Rev. Lett. 85 (2000) 499–502, http://dx.doi.org/10.1103/PhysRevLett.85.499, arXiv:hep-th/0003118, [hep-th].
[41] B.J. Carr, K. Kohri, Y. Sendouda, J. Yokoyama, New cosmological constraints on primordial black holes, Phys. Rev. D 81 (2010) 104019, http://dx.doi.org/10.1103/PhysRevD.81.104019, arXiv:0912.5297, [astro-ph.CO].
[42] S. Clark, B. Dutta, Y. Gao, L.E. Strigari, S. Watson, Planck constraint on relic primordial black holes, Phys. Rev. D 95 (8) (2017) 083006, http://dx.doi.org/10.1103/PhysRevD.95.083006, arXiv:1612.07738, [astro-ph.CO].
[43] W. DeRocco, P.W. Graham, Constraining primordial black hole abundance with the Galactic 511 keV line, Phys. Rev. Lett. 123 (25) (2019) 251102, http://dx.doi.org/10.1103/PhysRevLett.123.251102, arXiv:1906.07740, [astro-ph.CO].
[44] R. Laha, Primordial black holes as a dark matter candidate are severely constrained by the Galactic Center 511 keV γ-ray line, Phys. Rev. Lett. 123 (25) (2019) 251101, http://dx.doi.org/10.1103/PhysRevLett.123.251101, arXiv:1906.09994, [astro-ph.HE].
[45] B. Dasgupta, R. Laha, A. Ray, Neutrino and positron constraints on spinning primordial black hole dark matter, Phys. Rev. Lett. 125 (2020) 101101, http://dx.doi.org/10.1103/PhysRevLett.125.101101, arXiv:1912.01014, [hep-ph].
[46] C. Keith, D. Hooper, 511 keV excess and primordial black holes, Phys. Rev. D 104 (6) (2021) 063033, http://dx.doi.org/10.1103/PhysRevD.104.063033, arXiv:2103.08611, [astro-ph.CO].
[47] S. Mittal, A. Ray, G. Kulkarni, B. Dasgupta, Constraining primordial black holes as dark matter using the global 21-cm signal with X-ray heating and excess radio background, JCAP 03 (2022) 030, http://dx.doi.org/10.1088/1475-7516/2022/03/030, arXiv:2107.02190, [astro-ph.CO].
[48] A.K. Saha, R. Laha, Sensitivities on nonspinning and spinning primordial black hole dark matter with global 21-cm troughs, Phys. Rev. D 105 (10) (2022) 103026, http://dx.doi.org/10.1103/PhysRevD.105.103026, arXiv:2112.10794, [astro-ph.CO].
[49] A.K. Saha, A. Singh, P. Parashari, R. Laha, Hunting primordial black hole dark matter in Lyman-α forest, arXiv:2409.10617, [astro-ph.CO].
[50] R. Laha, J.B. Muñoz, T.R. Slatyer, Integral constraints on primordial black holes and particle dark matter, Phys. Rev. D 101 (12) (2020) 123514, http://dx.doi.org/10.1103/PhysRevD.101.123514, arXiv:2004.00627, [astro-ph.CO].
[51] J. Berteaud, F. Calore, J. Iguaz, P.D. Serpico, T. Siegert, Strong constraints on primordial black hole dark matter from 16 years of integral/spi observations, Phys. Rev. D 106 (2) (2022) 023030, http://dx.doi.org/10.1103/PhysRevD.106.023030, arXiv:2202.07483, [astro-ph.HE].
[52] M. Korwar, S. Profumo, Updated constraints on primordial black hole evaporation, JCAP 05 (2023) 054, http://dx.doi.org/10.1088/1475-7516/2023/05/054, arXiv:2302.04408, [hep-ph].
[53] N.J.L.S. Lobos, A.M. Ncube, R.C. Pantig, A.S. Cornell, Analyzing the effect of higher dimensions on the black hole silhouette, deflection angles, and PINN approximated quasinormal modes, arXiv:2406.08078, [gr-qc].
[54] R.C. Hickox, R. Narayan, T.R. Kallman, Origin of the soft excess in x-ray pulsars, Astrophys. J. 614 (2004) 881–896, http://dx.doi.org/10.1086/423928, arXiv:astro-ph/0407115, [astro-ph].
[55] R. Takahashi, T. Nakamura, Wave effects in gravitational lensing of gravitational waves from chirping binaries, Astrophys. J. 595 (2003) 1039–1051, http://dx.doi.org/10.1086/377430, arXiv:astro-ph/0305055, [astro-ph].
[56] M. Tamta, N. Raj, P. Sharma, Breaking into the window of primordial black hole dark matter with x-ray microlensing, arXiv:2405.20365, [astro-ph.HE].
[57] V. De Luca, G. Franciolini, A. Riotto, On the primordial black hole mass function for broad spectra, Phys. Lett. B 807 (2020) 135550, http://dx.doi.org/10.1016/j.physletb.2020.135550, arXiv:2001.04371, [astro-ph.CO].
[58] X. Barcons, et al., [Athena team], Athena: the X-ray observatory to study the hot and energetic universe, J. Phys. Conf. Ser. 610 (1) (2015) 012008, http://dx.doi.org/10.1088/1742-6596/610/1/012008.
[59] F. Özel others, [Lynx Team], The Lynx mission concept study interim report, arXiv:1809.09642, [astro-ph.IM].
[60] S.N. Zhang, et al., [eXTP], eXTP – enhanced X-ray timing and polarimetry mission, Proc. SPIE Int. Soc. Opt. Eng. 9905 (2016) 99051Q, http://dx.doi.org/10.1117/12.2232034, arXiv:1607.08823, [astro-ph.IM].
[61] P.S. Ray others, [STROBE-X Science Working Group], STROBE-X: X-ray timing and spectroscopy on dynamical timescales from microseconds to years, arXiv:1903.03035, [astro-ph.IM].
[62] R. Ruffini, J.A. Wheeler, Introducing the black hole, Phys. Today 24 (1) (1971) 30, http://dx.doi.org/10.1063/1.3022513.
[63] R.P. Kerr, Gravitational field of a spinning mass as an example of algebraically special metrics, Phys. Rev. Lett. 11 (1963) 237–238, http://dx.doi.org/10.1103/PhysRevLett.11.237.
[64] E.T. Newman, R. Couch, K. Chinnapared, A. Exton, A. Prakash, R. Torrence, Metric of a rotating, charged mass, J. Math. Phys. 6 (1965) 918–919, http://dx.doi.org/10.1063/1.1704351.
[65] R. Wald, Gedanken experiments to destroy a black hole, Annals Phys. 82 (2) (1974) 548–556, http://dx.doi.org/10.1016/0003-4916(74)90125-0.
[66] R. Penrose, Gravitational collapse and space–time singularities, Phys. Rev. Lett. 14 (1965) 57–59, http://dx.doi.org/10.1103/PhysRevLett.14.57.
[67] J.A. de Freitas Pacheco, E. Kiritsis, M. Lucca, J. Silk, Quasiextremal primordial black holes are a viable dark matter candidate, Phys. Rev. D 107 (12) (2023) 123525, http://dx.doi.org/10.1103/PhysRevD.107.123525, arXiv:2301.13215, [astro-ph.CO].
[68] L.A. Anchordoqui, I. Antoniadis, D. Lüst, Dark dimension, the swampland, and the dark matter fraction composed of primordial near-extremal black holes, Phys. Rev. D 109 (9) (2024) 095008, http://dx.doi.org/10.1103/PhysRevD.109.095008, arXiv:2401.09087, [hep-th].
[69] J.S. Schwinger, On gauge invariance and vacuum polarization, Phys. Rev. 82 (1951) 664–679, http://dx.doi.org/10.1103/PhysRev.82.664.
[70] G.W. Gibbons, Vacuum polarization and the spontaneous loss of charge by black holes, Commun. Math. Phys. 44 (1975) 245–264, http://dx.doi.org/10.1007/BF01609829.
[71] W.A. Hiscock, L.D. Weems, Evolution of charged evaporating black holes, Phys. Rev. D 41 (1990) 1142, http://dx.doi.org/10.1103/PhysRevD.41.1142.
[72] D.N. Page, Particle emission rates from a Black Hole II: Massless particles from a rotating hole, Phys. Rev. D 14 (1976) 3260–3273, http://dx.doi.org/10.1103/PhysRevD.14.3260.
[73] Y. Bai, N. Orlofsky, Primordial extremal black holes as dark matter, Phys. Rev. D 101 (5) (2020) 055006, http://dx.doi.org/10.1103/PhysRevD.101.055006, arXiv:1906.04858, [hep-ph].
[74] N. Arkani-Hamed, L. Motl, A. Nicolis, C. Vafa, The string landscape, black holes and gravity as the weakest force, JHEP 06 (2007) 060, http://dx.doi.org/10.1088/1126-6708/2007/06/060, arXiv:hep-th/0601001, [hep-th].
[75] D. Harlow, B. Heidenreich, M. Reece, T. Rudelius, Weak gravity conjecture, Rev. Mod. Phys. 95 (3) (2023) 3, http://dx.doi.org/10.1103/RevModPhys.95.035003, arXiv:2201.08380, [hep-th].
[76] E.F. Eiroa, G.E. Romero, D.F. Torres, Reissner–nordstrom black hole lensing, Phys. Rev. D 66 (2002) 024010, http://dx.doi.org/10.1103/PhysRevD.66.024010, arXiv:gr-qc/0203049, [gr-qc].
[77] M. Calzà, D. Pedrotti, S. Vagnozzi, Primordial regular black holes as all the dark matter (I): tr-symmetric metrics, arXiv:2409.02804, [gr-qc].

[78] M. Calzà, D. Pedrotti, S. Vagnozzi, Primordial regular black holes as all the dark matter (II): non-tr-symmetric and loop quantum gravity-inspired metrics, arXiv:2409.02807, [gr-qc].
[79] G. Dvali, A microscopic model of holography: Survival by the burden of memory, arXiv:1810.02336, [hep-th].
[80] G. Dvali, L. Eisemann, M. Michel, S. Zell, Black hole metamorphosis and stabilization by memory burden, Phys. Rev. D 102 (10) (2020) 103523, http://dx.doi.org/10.1103/PhysRevD.102.103523, arXiv:2006.00011, [hep-th].
[81] A. Alexandre, G. Dvali, E. Koutsangelas, New mass window for primordial black holes as dark matter from memory burden effect, arXiv:2402.14069, [hep-ph].
[82] G. Dvali, J.S. Valbuena-Bermúdez, M. Zantedeschi, Memory burden effect in black holes and solitons: Implications for PBH, arXiv:2405.13117, [hep-th].
[83] V. Thoss, A. Burkert, K. Kohri, Breakdown of Hawking evaporation opens new mass window for primordial black holes as dark matter candidate, arXiv:2402.17823, [astro-ph.CO].
[84] S. Balaji, G. Domènech, G. Franciolini, A. Ganz, J. Tränkle, Probing modified Hawking evaporation with gravitational waves from the primordial black hole dominated universe, arXiv:2403.14309, [gr-qc].