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Greybody factor and power spectra of the Hawking radiation in the 4*D* Einstein–Gauss–Bonnet de-Sitter gravity

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Abstract A novel 4D Einstein–Gauss–Bonnet gravity was recently formulated by Glavan and Lin [Phys. Rev. Lett. 124, 081301 (2020)]. Although this theory may run into trouble at the level of action or equations of motion, the spherically symmetric black hole solution, which can be successfully reproduced in those consistent theories of 4D EGB gravity, is still meaningful and worthy of study. In this paper, we investigate Hawking radiation in the spacetime containing such a de Sitter black hole. Both the greybody factor and the power spectra of the Hawking radiation of the massless scalar are studied numerically for the full range of various parameters, including the GB coupling constant α, the cosmological constant Λ and the coupling constant related to the scalar filed ξ. In particular, we find a negative α leads to a larger greybody factor than that of a $\alpha \geq 0$. While, for the power spectra of the Hawking radiation the situation is quite the opposite. The reason is that the temperature of the black hole would be very high when $\alpha < 0$. Actually, we observe that the temperature would be arbitrarily high when α approaches to the lower bound.

1 Introduction

Hawking radiation is one of the most important discoveries in quantum gravity, although it was derived semi-classically [1] in virtue of quantum field theory in cured spacetime. Since then, people believed black holes do have temperatures and satisfy the laws of thermodynamics [2]. On the other hand, Hawking radiation plays a central role in the study of the black hole information paradox, which may be the toughest obstacle in the way to understand quantum gravity thoroughly [3]. Based on the awareness of its importance and necessity, a lot of significant progresses have been made in the previous studies both experimentally and theoretically. On the experimental side, it was found that the Hawking radiation of tiny black holes may be observed by particle colliders through LHC [4–7], while Hawking radiation is too small to be directly detected through astronomical observations. Instead, Hawking radiation was proved to work in the laboratory analogues in terms of the theory of analogue gravity [8]. Further, some analogues of black-hole horizons in the laboratory have been realized and many evidences came into being to support the universality of Hawking radiation [9,10], see Ref. [11] for more complete review . On the purely theoretical side, Hawking radiation also has drawn attentions on many aspects of gravity, mainly including quantum information theory of black hole [12], holographic dual in the framework of AdS/CFT correspondence [13] and Hawking emission spectrum described in terms of the greybody factor in general relativity, extra-dimension models and other alternative gravity theories [14–43].

Thereinto, the Einstein–Gauss–Bonnet (EGB) gravity is known to be one of the most promising candidates for modified gravity [44]. A well-defined black hole solution for any dimension ($D \geq 5$) has been found [45] and many non-trivial effects indicate the study of EGB gravity not only helps us understand black holes and black branes more deeply [46–49] but also advances the development of the holographic gravity in the AdS/CFT correspondence [50]. Among these studies, many important results related to Hawking radiation in EGB gravity were also obtained and aroused a lot of concerns and discussions [19,20]. Despite the achievements, so far all the studies concerning the EGB gravity are limited to dimensions higher than four, so a connection with real black holes cannot be built.

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Recently, a novel 4D EGB gravity was formulated by rescaling the GB coupling constant which completed the missing piece of the EGB gravity. What is even more exciting is a spherically symmetric black hole solution in this theory derived by Glavan and Lin [51] is free of the singularity problem [52–54]. In fact, the same solution has been already found in other theories. For the first time it was discovered from the gravity with a conformal anomaly [52] and then it was recovered in the gravity with quantum corrections [53,54]. This 4D EGB gravity has aroused great interest since its publication [55–69]. Particularly, as many works have pointed out the procedure of taking $D \to 4$ limit in [51] may not be consistent [70–76]. On the other hand, some proposals have been raised to circumvent the issues of the novel 4D EGB gravity, including adding an extra degree of freedom to the theory [55,56,77,78] or breaking the temporal diffeomorphism invariance [79]. While, the novel $4D$ EGB gravity formulated in [51] may run into trouble at the level of action or equations of motion, the spherically symmetric black hole solution derived in [51] and in early literatures [52–54] can be successfully reproduced in those consistent theories of $4D$ EGB gravity, which means the spherically symmetric black hole solution itself is meaningful and worthy of study.

It's worth mentioning that, in [59], the authors first discussed the negative GB coupling constant case in details and found a negative GB coupling constant is allowed to retain a black hole solution, and some appealing features in this case were obtained. However, one doesn't know yet if a negative coupling constant still works when a positive cosmological constant is added in. Furthermore, the works on Hawking radiation are also absent.

In this paper, we focus on the 4D spherically symmetric EGB-dS black hole solution and investigate the full range of the GB coupling constant to retain a dS black hole. Because of the positive cosmological constant, we found a negative GB coupling constant is also allowed with both the GB coupling constant and cosmological constant are constrained non-trivially. After having the preparation, we perform a complete study on the greybody factor of the Hawking radiation firstly and proceed to the power spectra of the Hawking radiation of the massless scalar in the 4D GB-dS black hole background. The effects of various parameters, including the GB coupling constant, the cosmological constant as well as the coupling constant related to the scalar field, on both aspects of Hawking radiation are analyzed in detail. In particular, some new features are found when the GB coupling constant is negative.

The paper is organized as follows. In Sect. 2, we give a short review on the novel 4D EGB gravity and find the full range of the GB coupling constant and cosmological constant when the spacetime contains a dS black hole. In Sect. 3, we introduce the scalar field perturbation. Next, we turn our attention to the greybody factor in Sect. 4. And the energy emission rate of Hawking radiation are discussed in Sect. 5. We close our paper with a conclusion in Sect. 6.

2 The 4-dimensional EGB-dS black hole solution

The Einstein–Gauss–Bonnet gravity with a positive cosmological constant in D-dimensional spacetime has the action of the form

$$S_G = \frac{1}{16\pi G} \int d^D x \sqrt{-g}\, [R - 2\Lambda + \alpha \mathcal{L}_{GB}]\,, \tag{2.1}$$

where G is the D-dimensional Newton's constant, R is the Ricci scalar, Λ is the cosmological constant, and α is the Gauss–Bonnet (GB) coupling constant of dimension $(length)^2$. The Gauss–Bonnet term is given by

$$\mathcal{L}_{GB} = R_{\mu\nu\rho\sigma} R^{\mu\nu\rho\sigma} - 4 R_{\mu\nu} R^{\mu\nu} + R^2, \tag{2.2}$$

which is a total derivative in four dimensional spacetime and thus has no contribution to the dynamics in general. Surprisingly, a novel theory was discovered recently by rescaling the coupling constant as

$$\alpha \to \frac{\alpha}{D-4}. \tag{2.3}$$

and taking the limit $D \to 4$. Then the Lovelock's theorem is circumvented and a new spherically symmetric black hole solution was found [51].

The spherically symmetric GB-dS black hole solution of (2.1) in four dimensional spacetime can be described by

$$ds^2 = -f dt^2 + \frac{dr^2}{f} + r^2 d\Omega_2^2, \qquad f = 1 + \frac{r^2}{2\alpha}\left(1 - \sqrt{1 + \frac{4\alpha m}{r^3} + \frac{4\alpha\Lambda}{3}}\right). \tag{2.4}$$

Here $d\Omega_2^2$ is the line element of the 2-dimensional unit sphere S^2. The parameter m is related to the black hole mass M by $m = 2GM$. As we mentioned before, the above procedure to obtain the novel theory has been called into question by many follow-up works [70–76]. Whereas the above black hole solution (2.4) can be reproduced from those consistent $4D$ EGB gravity theories [55,56,77–79] and thus is exempted from being meaningless. In the following all our discussion will be based on the black hole solution (2.4), which can be independent of the original gravity theory [51].

This solution has the same form as its higher dimensional companions. But there are important differences. The four dimensional solution (2.4) can have three horizons in some parameter region, while the higher dimensional solutions have at most two horizons. The center singularity of (2.4) is repulsive, while it is attractive in higher dimensions. Thus the four dimensional solution is free from the singularity problem. Motivated by these differences, we study the Hawking radiation of the solution (2.4) in this paper.

We fix $r_h = 1$ for convenience, then the mass parameter m can be expressed in terms of α and Λ as

$$m = 1 + \alpha - \frac{\Lambda}{3} > 0. \tag{2.5}$$

As a consequence, the free parameters of the background are α, Λ now. In this paper, we focus on the parameter region where both the black hole event horizon and cosmological horizon exist. In addition, after plugging the Eq. (2.5) into $f(r)$, we find

$$f(\infty) = -\infty, \qquad f(0^+) = 1, \qquad f'(0^+) < 0 \tag{2.6}$$

for $\alpha > 0$, while for $\alpha < 0$, one has

$$f(\infty) = -\infty, \qquad 0 < r_0 < 1, \tag{2.7}$$

where r_0 is the root when the quadratic radical of the metric function $f(r)$ is vanishing and we have taken into account the condition $m > 0$ for the inequalities. By using these conditions in combination and following the similar analysis in [58], we conclude $f(1) = 0$ and $f'(1) > 0$ give the necessary and sufficient condition that this spacetime always contains a dS black hole, that is, both the event horizon and the cosmological horizon exist. Thus, we find the allowed region and show it in Fig. 1.

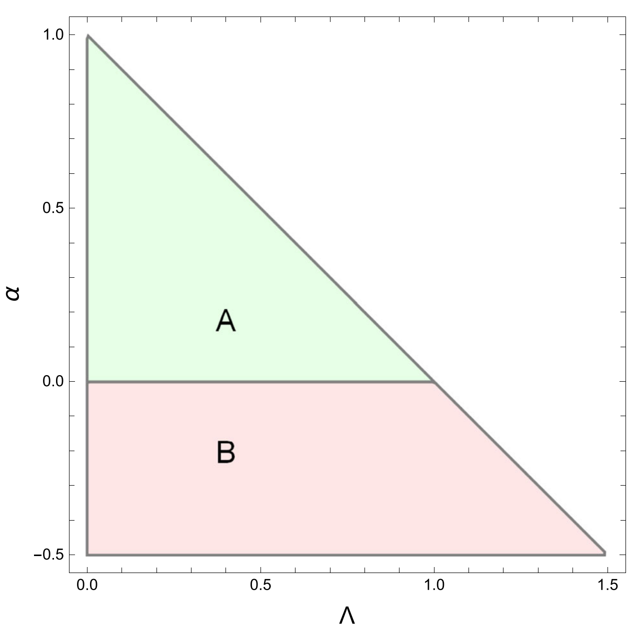

Fig. 1 The parameter region where both the event horizon and cosmological horizon exist. The region is bounded by $\alpha > -0.5$, $\Lambda > 0$ and $\alpha + \Lambda < 1$. The metric function is real for all $r > 0$ in region A, while becomes complex for small r in region B

3 The scalar field perturbation

We focus on the Hawking radiation of a massless scalar field Φ in the background (2.4). The scalar field Φ couples to gravity minimally or non-minimally with coupling constant ξ. The corresponding equation of motion is [25,26]

$$\nabla_\mu \nabla^\mu \Phi = \xi R \Phi. \tag{3.1}$$

For fluctuations of order $\mathcal{O}(\epsilon)$, the induced changes of the spacetime geometry are of the order $\mathcal{O}(\epsilon^2)$. Thus the effect of Φ on the background spacetime is negligible [43]. The above equation can be solved in a fixed background given by (2.4).

For stationary spherically symmetric background, the scalar wave function can be decomposed as

$$\Phi(t, r, \Omega) = \int d\omega \sum_{lm} e^{-i\omega t} \phi(r) Y_{lm}(\Omega), \tag{3.2}$$

where $Y_{lm}(\Omega)$ are spherical harmonics of the scalar wave function on S^2. The angular and radial part are decoupled and the radial master equation reads

$$\frac{1}{r^2}\frac{d}{dr}\left(f r^2 \frac{d\phi}{dr}\right) + \left[\frac{\omega^2}{f} - \frac{l(l+1)}{r^2} - \xi R\right]\phi = 0, \tag{3.3}$$

in which the Ricci scalar can be written as

$$R = -\partial_r^2 f + \frac{2}{r^2}\left(-2r\partial_r f + 1 - f\right). \tag{3.4}$$

The Eq. (3.3) can be written in the Schrödinger-like form

$$\frac{d^2 u}{dr_*^2} + (\omega^2 - V_{\text{eff}})u = 0, \tag{3.5}$$

where r_* is the tortoise coordinate defined by $dr_* = \frac{1}{f}dr$, and the new variable $u(r) = r\phi(r)$ is introduced. The effective potential

$$V_{\text{eff}} = f\left(\frac{l(l+1)}{r^2} + \xi R + \frac{1}{r}\partial_r f\right), \tag{3.6}$$

vanishes at the horizons of the spacetime and has a potential barrier between the event horizon r_h and the cosmological horizon r_c. The effective potential encodes the information of the background and the scalar field Φ, such as the angular momentum number l, scalar coupling constant ξ, cosmological constant Λ and GB coupling constant α. We will study the effects of these parameters on the Hawking radiation in detail in this paper.

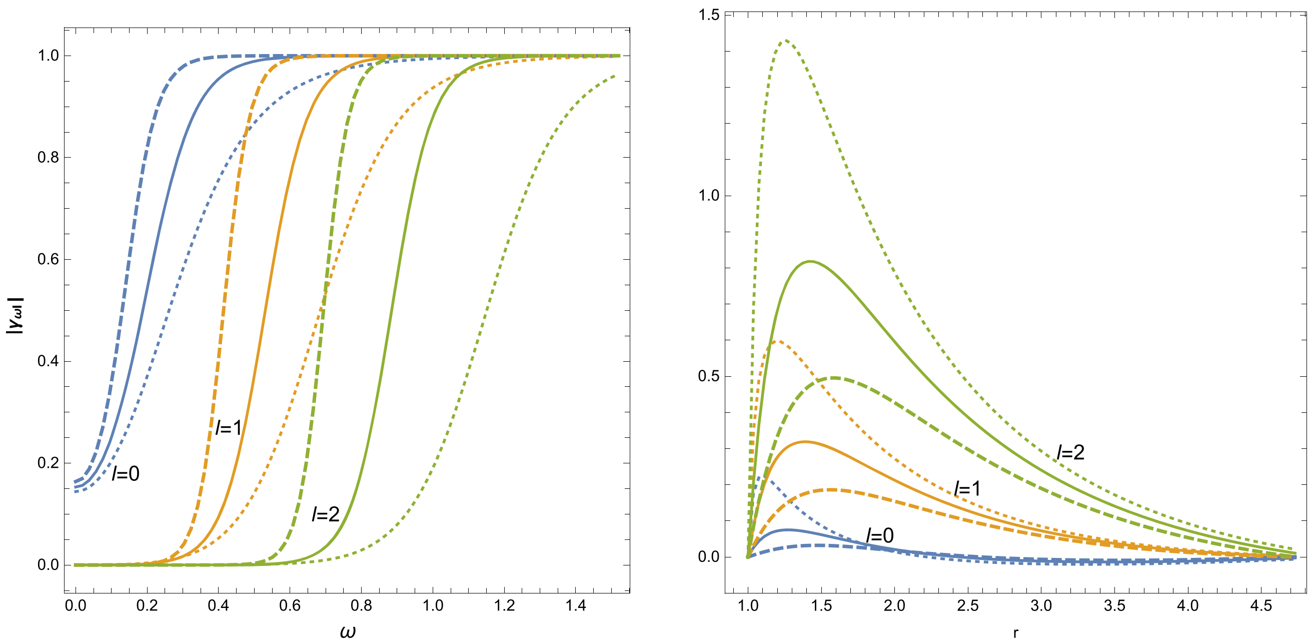

Fig. 2 Effects of α on the greybody factor for different modes l (left) and the corresponding effective potential (right). Dotted lines, solid lines and dashed lines denote $\alpha = -0.3$, $\alpha = 0$ and $\alpha = 0.3$, respectively. We fix $\Lambda = 0.1$ and $\xi = 0$ here

The radial equation (3.5) has asymptotic behavior

$$u \to \begin{cases} \mathcal{T}e^{-i\omega r_*} + \mathcal{O}e^{i\omega r_*} = \mathcal{T}(r-r_h)^{i\frac{\omega}{\kappa_h}} + \mathcal{O}(r-r_h)^{-i\frac{\omega}{\kappa_h}}, & r \to r_h, \\ \mathcal{R}e^{-i\omega r_*} + \mathcal{I}e^{i\omega r_*} = \mathcal{R}(r-r_c)^{i\frac{\omega}{\kappa_c}} + \mathcal{I}(r-r_c)^{-i\frac{\omega}{\kappa_c}}, & r \to r_c, \end{cases} \tag{3.7}$$

near the event horizon and cosmological horizon. Factors κ_h and κ_c are the surface gravity on r_h and r_c, respectively. The coefficients $\mathcal{I}$, $\mathcal{R}$ and $\mathcal{T}$ are the amplitudes of the incident, reflected and transmitted waves, respectively. The $\mathcal{O}$ term describes an outgoing wave at the event horizon and will be set to zero hereafter. Once these coefficients are worked out, we can get the greybody factor of the Hawking radiation by the definition

$$|\gamma_{\omega l}|^2 = 1 - |\mathcal{R}|^2/|\mathcal{I}|^2. \tag{3.8}$$

In general, the radial equation is hard to solve analytically. We thus turn to solve the radial equation numerically. We impose the ingoing boundary condition near the event horizon r_h and integrate the radial equation (3.3) towards the cosmological horizon r_c. By comparing with the asymptotic behavior (3.7) near r_c, we can get the coefficients $\mathcal{R}$, $\mathcal{I}$ and therefore the greybody factor (3.8). There exist various numerical methods to solve the radial equation. Here we adopt the method developed in [23–26].

4 The greybody factor of Hawking radiation

The free parameters appearing in (3.8) are α, Λ and ξ, l. We study their effects on the greybody factor in detail in this section.

4.1 Effects of α and l on the greybody factor

When $\alpha \to 0$, the solution (2.4) reduces to the Schwarzschild-dS (SdS) black hole. As α increases, the geometry is changed significantly due to the appearance of an inner horizon in the black hole. One may expect that the greybody factor of the Hawking radiation will also be affected significantly.

We show the effects of α on the greybody factor for different modes l in the left panel of Fig. 2. The greybody factor is enhanced for all modes l as α increases, and is suppressed by l when the other parameters are fixed. This phenomenon can be understood intuitively from the viewpoint of the effective potential, which is shown in the right panel of Fig. 2. The larger l, the higher potential barrier, so the wave is more likely to be reflected, which leads to a smaller greybody factor according to (3.8). On the other hand, the potential barrier decreases with α when the other parameters are fixed. Then the wave is unlikely to be reflected, which leads to a larger greybody factor. In fact, it can be shown numerically that the maximum of the effective potential increases monotonically with l and decreases monotonically with α. Thus l always suppresses the greybody factor and α always enhances it.

We would like to stress that a positive α makes the peak of the effective potential smaller, while a negative α increases the peak of the effective potential, which implies the existence of a negative α would make it harder for Hawking radiation to get through the barrier. This is a new phenomenon for which one may think the black holes with a negative GB coupling constant are harder to evaporate from the point of view of the greybody factor. But this is not true since the proper way to characterize the evaporation of black holes is

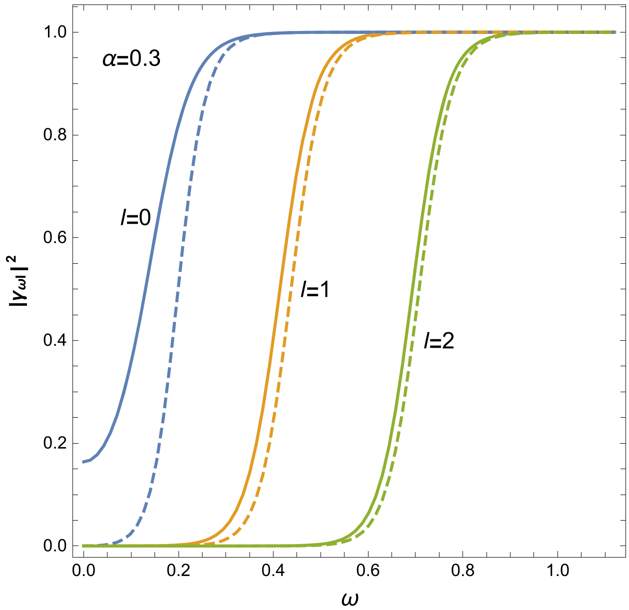

Fig. 3 Effects of ξ on the greybody factor for different l. Solid lines for $\xi = 0$ and dashed lines for $\xi = 0.3$. We fix $\Lambda = 0.1$ and $\alpha = 0.3$ here

via the energy emission rate which will be investigated in the next section.

Note that for the dominant mode $l = 0$, the greybody factor does not vanish when $\omega \to 0$. This is a characteristic feature of the minimally coupled massless scalar field propagating in the dS spacetime [16,18]. We can see that the presence of a GB term does not change this picture qualitatively. In the following subsection, we will see that this feature is changed qualitatively for nonminimally coupled scalar field.

4.2 Effect of ξ on the greybody factor

The effects of ξ on the greybody factor are shown in Fig. 3. We see that ξ decreases the greybody factor when the other parameters are fixed. Similarly, this phenomenon can also be understood intuitively from the effective potential. Now we know that a positive α increases the greybody factor, while ξ decreases it. There must be competition between α and ξ. However, we find that the effect of α is much weaker than that of ξ. In the presence of α, the derivative of the maximum of the effective potential with respect to ξ depends on the sign of curvature R, as can be seen from (3.6). To increase the greybody factor, this derivative must be negative which is possible only when $\Lambda < 3/56$. Above this threshold, the greybody factor always decreases with ξ in the whole allowed range of α. This phenomenon exists also in higher dimensional spacetimes. For example, the threshold is $\Lambda < 1/7$ when $D = 5$ and $\Lambda < 40/143$ when $D = 6$. The fact that the threshold is larger in higher dimensions implies the effect of α gets stronger in higher dimensions.

Another interesting property of ξ is that it introduces an effective mass for the scalar field, according to (3.1). Thus for the dominant mode $l = 0$ with frequency $\omega = 0$, the scalar field cannot pass through the effective potential barrier, so it leads to a vanishing greybody factor. This can be seen in Fig. 3.

4.3 Effect of Λ on the greybody factor

In the previous subsection, we see that the cosmological constant plays a subtle role for greybody factor. We study its effect on the greybody factor in detail in this subsection. From Fig. 4, we see that for positive α, when ξ is small, Λ increases the greybody factor. When ξ is large, Λ decreases the greybody factor. Similar behavior has been observed in higher dimensions [23–26]. On the other hand, Λ can be understood as an homogeneously distributed energy in the spacetime and boosts the particles to cross the effective potential barrier. However, it also contributes to the effective mass through (3.1) and so hinders the particles from crossing the barrier. The first effect dominates when ξ is small and the

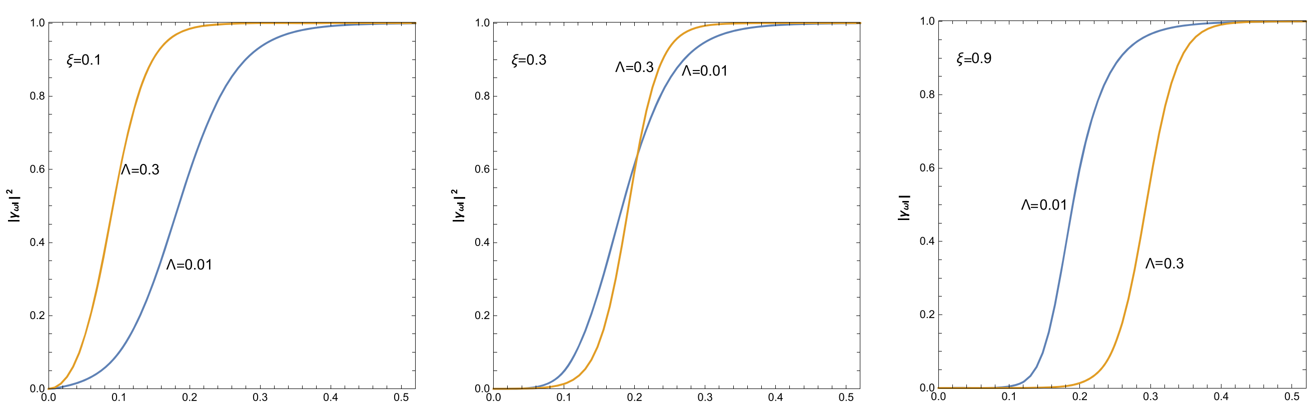

Fig. 4 Effect of Λ on the greybody factor. We fix $l = 0, \alpha = 0.3$ here

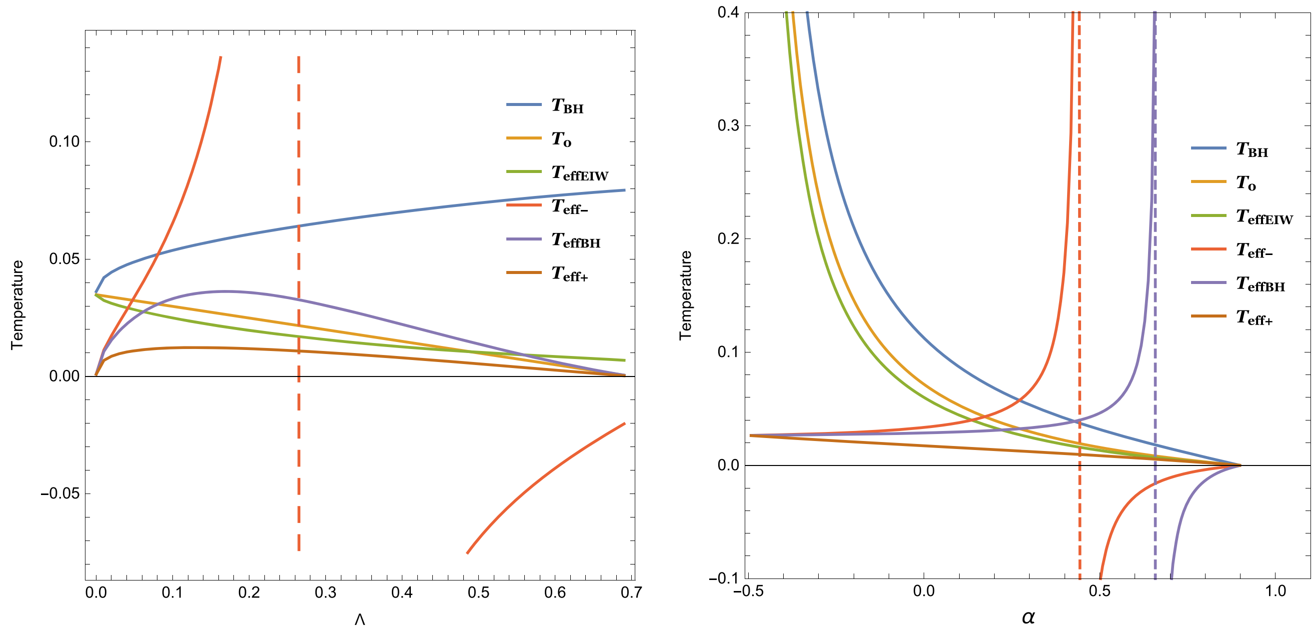

Fig. 5 The dependence of the effective temperatures of the black hole (2.4) on Λ (left, with fixed $\alpha = 0.3$) and α (right, with fixed $\Lambda = 0.1$). The extremal black hole satisfies $\alpha + \Lambda = 1$ and $\alpha > -0.5$

second effect dominates when ξ is large. When α is negative enough, we find that Λ increases the greybody factor in almost the whole frequency region. The effects of Λ on the effective mass is always weaker.

5 The power spectra of Hawking radiation

Once the greybody factor is worked out, the power spectra for Hawking radiation can be obtained by definition [17,18]

$$\frac{d^2E}{dt d\omega} = \frac{1}{2\pi}\sum_l \frac{N_l |\gamma_{\omega l}|^2 \omega}{e^{\omega/T} - 1}, \tag{5.1}$$

where T is the temperature of the black hole and $N_l = \frac{(2l+d-3)(l+d-4)!}{l!(d-3)!}$ is the multiplicity of the states that have the same l. As we have learned in the last section that the greybody factor of higher mode l is non-vanishing only when the frequency is high enough. While the power spectral is suppressed exponentially at high frequency according to (5.1). Thus only the lower modes l contribute to the total power spectra significantly. We take $l \leq 6$ for the numerical calculations hereafter.

For black holes in dS spacetime, the temperature of the system is subtle. One can define the temperature $T_0 = \frac{\kappa_h}{2\pi}$ on the event horizon and $T_c = -\frac{\kappa_c}{2\pi}$ on the cosmological horizon. T_0 is different with T_c in general such that the system is not in equilibrium. Inspired from the black hole thermodynamics, various effective temperatures were proposed [24,31,32], such as $T_{\text{eff-}} = \left(\frac{1}{T_c} - \frac{1}{T_0}\right)^{-1}$, $T_{\text{eff+}} = \left(\frac{1}{T_c} + \frac{1}{T_0}\right)^{-1}$, $T_{\text{effEIW}} = \frac{r_h^4 T_c + r_c^4 T_0}{(r_h+r_c)(r_c^3 - r_h^3)}$ and $T_{\text{BH}} = \frac{T_0}{\sqrt{f(r_0)}}$, $T_{\text{effBH}} = \left(\frac{1}{T_c} - \frac{1}{T_{BH}}\right)^{-1}$. Here r_0 is position where $\partial_r f|_{r=r_0} = 0$. Around this position, the black hole attraction balances the cosmological repulsion.

The dependence of these effective temperatures on Λ is shown in the left panel of Fig. 5. We see that T_{BH} increases with Λ, while T_0 and T_{effEIW} decrease with Λ monotonically. T_{effBH} and $T_{\text{eff+}}$ have a maximum. $T_{\text{eff-}}$ diverges at $\Lambda = 0.265$. The dependence of the effective temperatures on α is shown in the right panel of Fig. 5. Very interestingly, we find the three effective temperatures go to T_c when $\alpha = -0.5$. The reason is that the temperature of the event horizon of the black hole T_{BH} and the temperature defined on the event horizon T_0 approach to infinity for $\alpha = -0.5$, that is to say, the black hole is unlimited hot for this case which is worthy of further study. Also, we see that T_{BH}, T_0, T_{effEIW} and $T_{\text{eff+}}$ decrease with α. It's worth taking a moment to flag the fact all the temperatures are continuous across $\alpha = 0$ although the black hole for a positive and negative α seem different. In addition, we find $T_{\text{eff-}}$ diverges at $\alpha = 0.443$ and T_{effBH} diverges at $\alpha = 0.668$. To get reasonable result for the power spectra, we should abandon $T_{\text{eff-}}$ and T_{effBH}. For both small Λ and large Λ, $T_{\text{eff+}}$ tends to zero, which will lead to vanishing power spectra according to (5.1). This is unreasonable and thus we also abandon $T_{\text{eff+}}$. It has been shown in higher dimensions that only T_{BH} leads to significant radiation in the whole parameter space [26]. In this section,

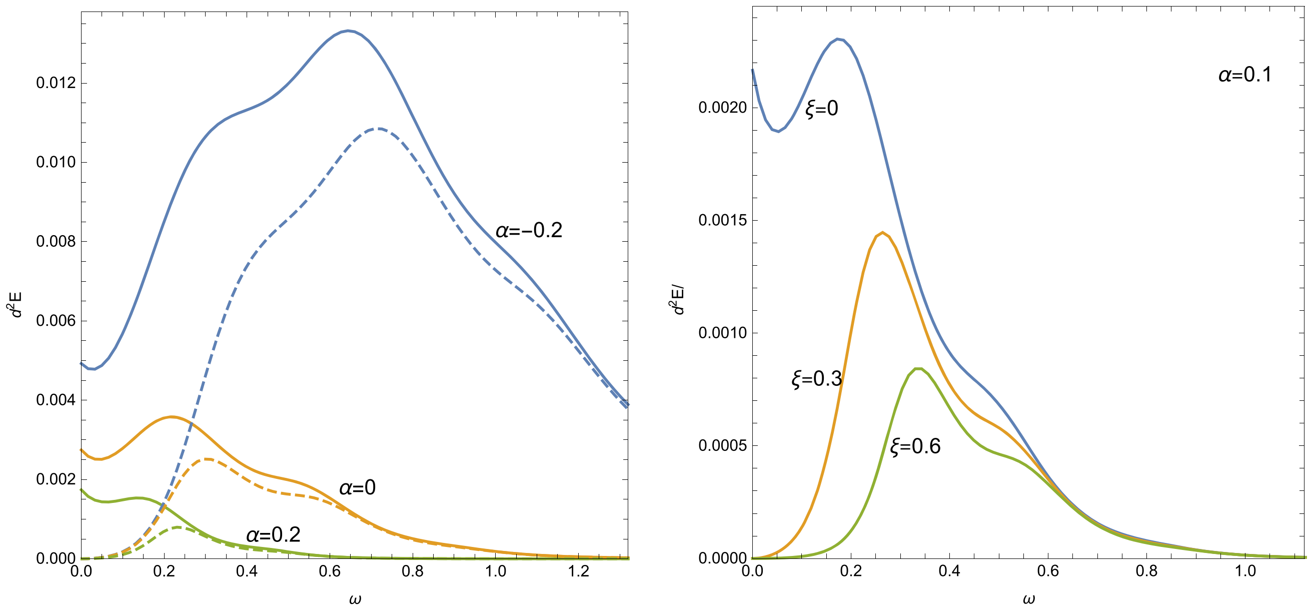

Fig. 6 The effects of α (left panel, solid lines for $\xi = 0$, dashed lines for $\xi = 0.3$) and ξ (right panel, with fixed $\alpha = 0.1$) on the power spectra of Hawking radiation

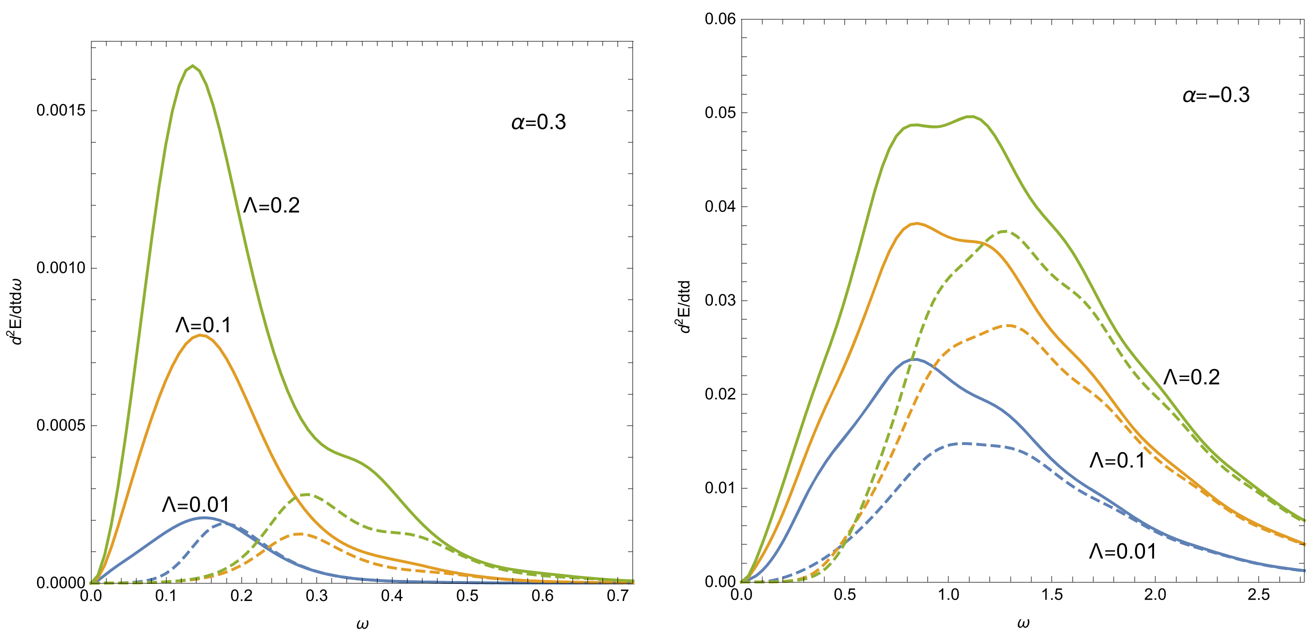

Fig. 7 Effect of Λ on the power spectra of Hawking radiation. Solid lines for $\xi = 0.1$, dashed lines for $\xi = 0.6$

we therefore take T_{BH} to study the power spectra of Hawking radiation.

5.1 Effects of α and ξ on the power spectra

The effects of α on the power spectra of Hawking radiation are shown in the left panel of Fig. 6. For both minimally and nonminimally coupled scalar field, the power spectra is suppressed as α increases, which is contrast with its effect on the greybody factor. In particular, we find that line of the power spectra with $\alpha < 0$ is above other lines with $\alpha \geq 0$. In other words, the intensity of Hawking radiation is much high with a negative α, although the greybody factor is lower. The reason comes from that the temperature of the black hole is very high when α becomes negative. In fact, the temperature plays a more important role here. T_{BH} decreases with α, and the power spectra also decreases with α according to (5.1). Note that for minimally coupled scalar field, the power spectra at low frequency are non-vanishing due to the finite greybody factor there. For nonminimally coupled scalar field, the power spectra at low frequency tends to zero. The peak of the power spectra moves to lower frequency as α increases, which is consistent with Fig. 4.2 where the greybody factor moves to the left as α increases.

The effects of ξ on the power spectra are shown in the right panel of Fig. 6. The power spectra are suppressed by ξ in the whole frequency region. We have learned that in Sect. 4.2, the greybody factor can be enhanced at high frequency when $\Lambda < 3/56$. However, (5.1) tells us that the power spectra are exponentially suppressed in high frequency. Thus the qualitative behavior that ξ suppresses the power spectra is independent of α. Note that the peak of the power spectra for nonminimally coupled scalar field moves to higher frequency when ξ increases. This is consistent with Fig. 3 where the greybody factor moves to the right as ξ increases.

5.2 Effect of Λ on the power spectra

The effect of Λ on the power spectra of the Hawking radiation is more subtle. For positive α, as shown in the left panel of Fig. 7, the cosmological constant Λ enhances the power spectra in the whole frequency region when ξ is small. When ξ is large, Λ enhances the power spectra only in the high frequency region. In the low frequency region, it suppresses the power spectra, as shown by the dashed lines in Fig. 7. This behavior is consistent with Fig. 4, where the cosmological constant enhances the greybody factor when ξ is small and suppresses the greybody factor when ξ is large. However, when α is negative enough, the cosmological constant enhances the power spectra in almost the whole frequency region no matter how large ξ is, as shown in the right panel of Fig. 7.

6 Summary

The four dimensional Einstein–Gauss–Bonnet black holes found recently have some distinct properties compared to their higher dimensional companions. For example, the spherically symmetric neutral 4D black hole solutions in asymptotically de Sitter spacetime can have three horizons, while the solutions in higher dimensional spacetime have only two horizons. On the other hand, we also found the spacetime contains a black hole when the GB coupling constant is negative. Furthermore, the EGB black hole in 4D can be directly compared with the Schwarzschild black hole which is the most common model used to describe a real black hole in our universe. At this point, 4D EGB black hole has a significant advantage over the GB black holes in higher dimensions. One thus expect that the Hawking radiation may also have distinct properties compared to the higher dimensional case.

We studied the greybody factor of the Hawking radiation firstly. The greybody factor is suppressed heavily by the angular momentum number l of the scalar mode. The Gauss–Bonnet coupling constant enhances the greybody factor, while the nonminimally coupling constant ξ of the scalar field decreases it. In particular, we found compared to Schwarzschild-dS black hole ($\alpha = 0$), the 4D EGB black hole with a negative α has a larger greybody factor. When the frequency of the mode tends to zero, the greybody factor vanishes for nonminimally coupled massless scalar field while has a finite value for minimally couple massless scalar field. The role of cosmological constant in greybody factor depends on ξ. It enhances the greybody factor when ξ is small, and decreases it when ξ is large. These behaviors are similar to the higher dimensional cases qualitatively, and can be understood intuitively from the viewpoint of the effective potential.

We then studied the power spectra of the Hawking radiation of the massless scalar field in the 4D GB-dS background. We analysed various definitions of temperature in asymptotically dS spacetime and adopted the most reasonable one to calculate the power spectra of the Hawking radiation. We found that both ξ and α suppress the power spectra. In particular, the power spectra with a negative α is over those with $\alpha \geq 0$, which means the intensity of Hawking radiation with $\alpha < 0$ is higher than others with $\alpha \geq 0$, although the greybody factor is opposite. The cosmological constant enhances the power spectra when ξ is small and suppresses it when ξ is large. The power spectra vanish for nonminimally coupled scalar field while have a finite value for minimally couples scalar field.

The method used in this work can be developed to study the amplitude of the superradiance of the charged black holes in 4D Einstein–Gauss–Bonnet gravity. It is known that RN-dS black holes are linearly unstable to spherical charged scalar perturbations [42]. It can be expected that the charged 4D GB-(A)dS black holes have also superradiant instability.

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