# Hawking radiation via tachyon condensation and its implications to tachyon cosmology

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**Abstract**

Hawking radiation can be derived from the collapsing process of matter to form a black hole. In this work, we show in more detail that the freely infalling process of a probe (D-)particle (or point-like object) in a non-extreme black hole background is essentially a tachyon condensation process. That is, a probe D-particle will behave as an unstable D-particle in the near-horizon region of a non-extreme black hole. From this point of view, Hawking radiation can be viewed as the thermal radiation from rolling tachyon on an unstable D-particle (i.e., the infalling probe) at the Hagedorn temperature. The result has interesting implications to tachyon cosmology: the uniform tachyon rolling in cosmology can automatically create particle pairs at late times, via a mechanism just like the Hawking radiation process near a black hole. So this particle creation process can naturally give rise to a hot universe with thermal perturbations beyond tachyon inflation, providing an alternative reheating mechanism.

## 1 Introduction

It is known that strings can be created and radiated from evolving unstable D-branes or S-branes, whose dynamics can be described by various open-sting tachyon field theories that are believed to be equivalent. It is found in boundary conformal field theory (BCFT)

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in [1, 2, 3] that open-string pairs can be created from rolling tachyon at the Hagedorn temperature. This implies that most or all energy of the tachyon is transferred into closed string radiation [4, 5, 6, 7]. Indeed, the study [6, 7] shows that an unstable D-particle or wrapped unstable D$p$-brane ($p \leq 2$) will completely decay into closed strings, mainly massive ones. However, it is hard to extract exact information about this intense energy release process due to strong back reaction. Moreover, the results are mainly investigated in BCFT. In the tachyon effective field theory and other string field theories, they are not really seriously verified.

In a previous work [8], we found that a probe particle infalling in Rindler space should behave as an unstable particle. This reveals that the dynamics near non-extreme black holes can be described by a time-dependent tachyon field theory. This connection makes it possible to understand better one side of this dual-like description from the other. From knowledge of tachyon condensation, we learn that material collapsing into a non-extreme black hole should encounter a phase transition at the Hagedorn temperature and should decay into gravitons (or generally closed strings) completely, mainly very massive ones, before reaching the horizon. We have also shown that the information of the material collapsing into black holes should be stored in these gravitons.

On the other hand, one can also learn some useful information about tachyon condensation from black hole thermodynamics. First, we can know that the rolling tachyon should really produce thermal radiation, with the Hagedorn temperature equal to the corresponding Hawking temperature. Second, we may get more accurate information about the closed-string emission from tachyon condensation if we agree that the degenerated states of closed strings emitted from an infalling particle should be responsible for the increased entropy in the black hole due to absorption of the particle.

In this work, we shall focus on the correlation between Hawking radiation from black holes and thermal radiation from tachyon. We shall show, in more detail than in the previous work [8], that infalling particle attracted by gravitational force of a non-extreme black hole will behave as an unstable (D-)particle approaching the event horizon and will give rise to thermal particle-pair creation via tachyon condensation. We argue that this particle creation process from the infalling probe provides an alternative explanation to the Hawking radiation from the black hole. We further show that this has interesting applications in the tachyon cosmological model, which is also constructed based on the tachyon effective theory.

The tachyon driven cosmology [9, 10, 11, 12, 13] has the merit that it can give a unified description of inflation, dark matter and even dark energy [14] as time grows. However, this model is thought problematic in some aspects [12, 15, 13]. One of the key problems is that a sensible reheating mechanism expected after inflation is hard to be realised in this model since the tachyon field does not oscillate around the minimum of the asymptotic potential [13]. The mechanisms proposed in past years usually involve the assistance of companion fields: coupling the evolving tachyon to massless fields like gauge fields [16, 17, 18] or introducing extra curvaton field [19]. In this work, we show in terms of our above results that uniform tachyon rolling in cosmology can automatically create and radiate particles at late times after inflation, with no need of assistant fields. The tachyon produces Hawking-like radiation at the Hagedorn temperature, which is high in the early universe and cools down as the universe expands. This is supported by the perturbative analysis that indicates analogous behaviours in late-time tachyon cosmology and in a dS universe.

The paper is organised as follows. In Section. 2, we extend our previous suggestion that the infalling process towards a non-extreme black hole is essentially a tachyon condensation process. Based on the point, we argue that the Hawking radiation from black hole can be interpreted as thermal radiation from the infalling probe, which can be viewed as an unstable (D-)particle. We shall discuss their implications to tachyon cosmology in Section. 4, in assistance with the perturbative analysis given in Section. 5. In the last section, the conclusions are made. In the Appendix, we present the results about thermal radiation from rolling tachyon in BCFT and rederive them in the effective field theory.

## 2 The infalling process and tachyon condensation

In [8], we have shown that the dynamics of an infalling (D-)particle (or point-like object) in the near-horizon region (Rindler space) of a non-extreme black hole can be described by the tachyon effective theory, i.e., the probe (D-)particle in Rindler space behaves as an unstable (D-)particle. For an infalling particle of mass $\tau_0$ towards the horizon, it can be described by the tachyon effective action at large $|T|$:

$$S(T) = -\int dt V(T)\sqrt{1-\dot{T}^2}, \quad V(T) = \frac{\widetilde{\tau}_0}{\cosh(\beta T)}, \tag{1}$$

where $\widetilde{\tau}_0 = \kappa\tau_0$ (with $\kappa$ a constant) can be viewed as the mass of the unstable particle. This approach to yield an unstable D-particle configuration is similar to the method of constructing geometrical tachyon in [29, 30], where it was found that a prob D-brane on the geometry of wrapped NS5-branes behaves as an unstable D-brane. In what follows, we shall represent and extend the results.

## 2.1 Asymptotic region

Let us first extend to the asymptotic region of a non-extreme black hole satisfying asymptotically flat boundary condition. In this region, the geometry is nearly the Minkowski spacetime. It is easy to know that the dynamics of a probe on it can be described by the tachyon effective action (1) sitting on the top of the potential:

$$S_0 = -\widetilde{\tau}_0 \int dt \sqrt{1 - \dot{T}^2} = S(T \to 0). \tag{2}$$

## 2.2 Geometric tachyon in Rindler space

For non-extreme black holes, the geometry near the event horizon is generically the Rindler space. To describe the Rindler metric, let us start with the two-dimensional Minkowski spacetime as usual:

$$ds = -du^2 + dv^2. \tag{3}$$

(1) Right wedge: In terms of the trajectory of a uniformly accelerating observer, we do the following reparameterisations for the new spatial coordinate $T \geq 0$

$$u = \frac{2}{\beta} e^{-\beta T} \sinh(\beta t), \quad v = \frac{2}{\beta} e^{-\beta T} \cosh(\beta t), \tag{4}$$

so that $v \geq 0$. For $T = 0$, this is the trajectory of a uniformly accelerating observer with the acceleration of $\beta$. Note that the transformations here are a little bit different from those given in standard textbooks. For later convenience, we have chosen an overall factor 2 instead of 1 and a minus sign in the exponential function. But this does not make severe physical difference.

Under the transformations (4), we get the Rindler metric of the form on the right wedge

$$ds^2_{RR} = 4e^{-2\beta T}(-dt^2 + dT^2). \tag{5}$$

In regions closed to the event horizon, the coordinate $T$ is large. The action $S_0$ of a particle moving on the right Rindler wedge (5) is actually the effective action $S(T)$ (1) at large positive $T$

$$S_0 = -2\tau_0 \int dt e^{-\beta T} \sqrt{1 - \dot{T}^2} \simeq S\left(T \gg \beta^{-1}\right). \tag{6}$$

Now the coordinate $T$ finds the meaning of a scalar field in the action.

In this wedge, the time $t$ evolves from $-\infty$ to $\infty$. For large and constant $T$, the trajectories at large $|t|$ are the curves that are very closed to the null ones $u = \pm v$, where the horizons are located. In the tachyon field theory (6), the null curves correspond to the solutions at $|t| \to \infty$ and $T \to \infty$: $\dot{T} = \pm 1$.

(2) Left wedge: The reparameterisations on the $T \leq 0$ side are given by:

$$u = -\frac{2}{\beta} e^{\beta T} \sinh(\beta t), \quad v = -\frac{2}{\beta} e^{\beta T} \cosh(\beta t), \tag{7}$$

so that it covers the $v \leq 0$ part of the full Minkowski space. These lead to the left Rinlder wedge

$$ds^2_{LR} = 4e^{2\beta T}(-dt^2 + dT^2). \tag{8}$$

The particle in this wedge with large $-T$ is described by the effective action (1) on the negative $T$ side

$$S_0 = -2\tau_0 \int dt e^{\beta T} \sqrt{1 - \dot{T}^2} \simeq S\left(T \ll -\beta^{-1}\right). \tag{9}$$

In this wedge, we can assume that the time evolves from $\infty$ to $-\infty$. The trajectories for large and constant $-T$ are also those in the neighborhood of the null curves $u = \pm v$ (but on the opposite side).

The Rindler space can also be expressed as

$$ds^2_R = -\rho^2 dt^2 + \frac{1}{\beta^2} d\rho^2. \tag{10}$$

It has two branches $\rho \leq 0$ and $\rho \geq 0$, with the horizon located at $\rho = 0$. The coordinates can be related to those in (5) and (8) via the relations:

$$T = \mp \frac{1}{\beta} \ln\left(\pm \frac{\rho}{2}\right). \tag{11}$$

Thus, towards the horizon $\rho \to 0^{\pm}$, $T \to \pm\infty$ respectively.

However, the region in between the spatial infinity and the near-horizon Rindler wedges may not be described by the tachyon effective action (1). But this intermediate stage is usually physically unimportant.

For extremal black holes, the near-horizon geometry is AdS space instead of Rindler space. So the action of a probe particle does not take the form of the tachyon effective action. The infalling process should not be a tachyon condensation process.

## 2.3 Example

We can take a look at a simple and well-studied example, the Schwarzschild black hole:

$$ds^2 = -\left(1 - \frac{r_0}{r}\right) d\chi^2 + \left(1 - \frac{r_0}{r}\right)^{-1} dr^2 + r^2 d\Omega^2, \tag{12}$$

where $r_0 = 2M$. We redefine the coordinates

$$r = r_0(1 + \rho^2), \quad \chi = 2\beta r_0 t, \tag{13}$$

so that $\rho$ is dimensionless. So the first two components of the metric are

$$ds_2^2 = 4\beta^2 r_0^2 \left( -\frac{\rho^2}{1+\rho^2} dt^2 + \frac{1+\rho^2}{\beta^2} d\rho^2 \right). \tag{14}$$

Note that we now have two branches $\rho \geq 0$ and $\rho \leq 0$ (or $T \geq 0$ and $T \leq 0$ respectively as will be redefined later) even though $r$ is single valued.

At $r \to \infty$ or $\rho \to \pm\infty$, it is nearly flat spacetime. We may formally translate the radial coordinate $T = \mp[\rho^2/(2\beta) - c]$ by a constant $c \to \infty$, so that $T \to 0^\pm$ as $\rho \to \pm\infty$ correspondingly. So the action of a particle with mass $\tau_0$ moving along the radial direction of the metric (14) at spatial infinity is actually the the tachyon effective action $S(T)$ (1) on the top of the potential

$$S_0(\rho \to \pm\infty) = -2\beta r_0 \tau_0 \int dt \sqrt{1 - \dot{T}^2} \simeq 2\beta r_0 S(T \to 0^\pm). \tag{15}$$

The factor $\kappa = 2\beta r_0$ appears because $S(T)$ is defined in the time coordinate $t$ in stead of the original $\chi$, as present in Eq. (13). The mass of the particle measured in $\chi$ is $\tau_0$, while it is $2\beta r_0 \tau_0$ measured in $t$.

Approaching the horizon $r \to r_0$ or $\rho \to 0^\pm$, it is the Rindler space

$$ds_2^2 = 4\beta^2 r_0^2 \left( -\rho^2 dt^2 + \frac{1}{\beta^2} d\rho^2 \right). \tag{16}$$

The particle moving in the right and left Rinlder wedges are respectively

$$S_0(\rho \to 0^+) \simeq 2\beta r_0 S(T \to \infty), \tag{17}$$

$$S_0(\rho \to 0^-) \simeq 2\beta r_0 S(T \to -\infty), \tag{18}$$

via the redefinition relation (11). Thus, an ordinary particle in Rindler space can be viewed as an unstable particle.

In summary, the infalling process can be correlated with the tachyon condensation process, as presented in Table 1. The tachyon field theory on the top of the potential is located at the false vacuum, which is the open string vacuum. Approaching the horizon corresponds to the closed string vacuum $V \to 0$ of the theory, with no open strings present.

The above analysis can be extended to the extended Kruskal coordinates of Schwarzschild black holes. We can see that it is a tachyon condensation (brane decay) process collapsing from the past null infinity $\mathcal{I}^-$ to the black hole horizon $\mathcal{H}^+$, while it corresponds to the inverse process of tachyon condensation (brane creation) coming out from the white hole horizon $\mathcal{H}^-$ to the future null infinity $\mathcal{I}^+$. It is interesting to notice that it is a brane creation and then decay process from $\mathcal{H}^-$ to $\mathcal{H}^+$, i.e., the evolution of a full S-brane. All these process can happen in either the right or the left region.

| | Left ($\rho, T \leq 0$) | Right ($\rho, T \geq 0$) | tachyon potential |
|---|---|---|---|
| $r \to \infty$ ($\mathcal{I}^- \cup \mathcal{I}^+$) | $\rho \to -\infty$, $T \to 0^-$ | $\rho \to \infty$, $T \to 0^+$ | $V(T) \to \widetilde{\tau}_0$ |
| $r \to r_0$ ($\mathcal{H}^- \cup \mathcal{H}^+$) | $\rho \to 0^-$, $T \to -\infty$ | $\rho \to 0^+$, $T \to \infty$ | $V(T) \to 0$ |

Table 1: Correspondence between the infalling process and tachyon condensation at the two vacua.

# 3 Hawking radiation via tachyon condensation

It is known that particle (open-string) pairs of thermal distribution can be created at the critical Hagedorn temperature in tachyon field theory [1, 3] (see the Appendix).

Since the infalling process of a probe towards a non-extreme black hole is actually a tachyon condensation process, the Hawking radiation from the black hole could be interpreted as the thermal radiation from rolling tachyon on the unstable particle (i.e., the infalling probe). The source that makes the probe behave as an unstable particle is the gravitational force of the black hole. The energy of radiation from the unstable particle originally comes from the black hole. In what follows, we explore connections between the tachyon field theory approach and traditional approaches to Hawking radiation from black holes.

## 3.1 Unruh effects in Rindler space and in tachyon field theory

We first take a look at how the Unruh effect in Rindler space is relevant in the tachyon field theory, which actually describes an accelerating scalar field. The Unruh effect in Rindler space is usually analyzed by a probe scalar $\phi$ of mass $m_s$ (in some simple discussion, $m_s$ is set to be zero; the massive case is reviewed in [31] for example).

In the flat metric (3), the KG equation is

$$(\partial_u^2 - \partial_v^2 + m_s^2)\phi = 0, \tag{19}$$

whose solutions are plane waves. In Minkowski spacetime, it is convenient to define the positive frequency modes $f_k$ based on the solution $\phi$. This equation is similar to Eq. (54) (in the Appendix) around $T \simeq 0$ in the tachyon field theory if we ignore quadratic and higher order terms of $T$ with $\mathrm{sech}(\beta T) \simeq 1 - \mathcal{O}(T^2)$.

In the Rindler wedges (5) and (8), the KG equations are respectively

$$\left(\partial_t^2 - \partial_T^2 + m_s^2 e^{\mp 2\beta T}\right)\phi = 0, \tag{20}$$

as discussed in many references. This equation corresponds to Eq. (54) at large $|T|$ in the tachyon field theory. So the positive frequency modes $g_k^L$ and $g_k^R$ from the equation (see e.g. [31]) respectively corresponds to $e^{it}\psi^{(-)}$ and $e^{-it}\psi^{(+)}$ from Eq. (54) at large $|T|$. They have similar asymptotical forms to Eqs. (65) and (66) in the tachyon effective theory.

The expansions of $\phi$ in terms of the above modes are

$$\phi = \sum_k (a_k f_k + h.c.) = \sum_k (b_k g_k^L + c_k g_k^R + h.c.). \tag{21}$$

The Bogolubov coefficients in the transformations between $g_k^L$, $g_k^R$ and $f_k$ can be derived via the relations (4) and (7) between the coordinates $(u, v)$ and $(T, t)$, by extending the modes in Rindler wedges to the full Minkowski spacetime. But, in the tachyon field theory (as shown in the Appendix), the Bogolubov tansformations (67) and (68) are determined by the Legendre functions from a single continuum equation with a potential which smoothly connects the two vacua of the theory. But the temperatures of the thermal spectra obtained in the two approaches are the same. The Unruh effect in the BCFT has been discussed in [3].

### 3.2 Hawking radiation via particle creation from infalling probe

In Hawking's original work [32], thermal radiation from black hole was derived by considering a Schwarzschild black hole formed from collapsing matter. In our above discussion, we have argued that the infalling process of matter (which can be taken as a collection of point-like cells) can be viewed as a tachyon condensation process. So the Hawking radiation could be understood as the thermal radiation in the tachyon field theory.

Following Hawking's approach, we do not need to consider the white hole horizon. We only consider the collapsing process from $\mathcal{I}^-$ to $\mathcal{H}^+$, which have been argued above to correspond to the rolling process of tachyon from $T = 0$ to $T \to \infty$. For a propagating scalar through the collapsing system, normalise the outgoing modes at $\mathcal{I}^+$ as $\phi_{\text{out}} \sim e^{-i\omega(\chi - r^*)}$ and then track back in time to $\mathcal{I}^-$, so that we do not need to know the full analytical solution in all spacetime outside the horizon. Here, the tortoise coordinate $r^*$ is defined as $r^* = r + r_0 \ln(r/r_0 - 1)$ and is approximately $r^* \simeq -2\beta r_0 T$ near the horizon. With $\chi = 2\beta r_0 t$ in Eq. (14), the normalised outgoing mode takes the form $\phi_{\text{out}} \sim e^{-i2\beta r_0(t+T)}$ near the horizon in the Rindler coordinates $(T, t)$, which are defined in Eqs. (13) and (11). This is consistent with the normalised outgoing mode $e^{-it}\psi^{(+)} \sim e^{-i(t+T)}$ given in Eq. (65) in the tachyon field theory, where we have argued that this mode is chosen because it is relevant to the tachyon condensation process, but not the inverse process.

When we go back to $\mathcal{I}^-$, we can find the correlation of the outgoing mode with the ingoing mode $\phi_{\rm in} \sim e^{-i\omega(t+r^*)}$ at $\mathcal{I}^-$. This correlation gives rise to the Bogolubov coefficient that leads to a Planck spectrum of particle creation. The correlation can be determined in various methods, e.g., the geometric optics approach. Here, the method introduced in [33] is more relevant. Consider a collapsing thin shell, with light rays propagating through it during the collapsing process. In the interior of the shell, the metric is flat. Outside the shell, the metric is the Schwarzschild metric. The location of the shell $r = R(\chi)$ is determined by the junction condition on the boundary of the shell. When a light ray enter the shell at a distance finitely away from the horizon and then exit near the horizon $R \sim r_0$, the correlation and so the Bogolubov coefficient between the modes can be determined, which indeed lead to the thermal radiation spectrum of the background black hole. Thus, the Hawking radiation is derived by the aid of a probe shell.

Here, the situation is similar: the Hawking radiation can be derived by the aid of an infalling probe (D-)particle. As shown in the previous section, a probe (D-)particle in a non-extreme black hole spacetime will behave as an unstable (D-)particle in the near-horizon region. Particle pairs of thermal distribution at the Hagedorn temperature will be created from the probe during the collapsing process (i.e., the tachyon condensation process) (see the Appendix). This should be essentially the Hawking radiation of the black hole itself.

So the tachyon field theory approach to deriving Hawking radiation is compatible with previous approaches. The derived temperatures in these ways are the same. From this point of view, the Hawking temperature is actually the Hagedorn temperature in the tachyon field theory. They are equivalent when detected in the same time coordinate:

$$T^{(t)}_{\rm Haw} = \frac{\beta}{2\pi} = T^{(t)}_{\rm Hag}. \tag{22}$$

Detected in the Schwarzschild time coordinate $\chi$, the temperatures are both $1/(4\pi r_0)$ as expected, due to the coordinate relation (13).

This interpretation of Hawking temperature as the Hagedorn temperature means that the infalling particle will experience a Hagedorn phase transition, as discussed in the previous work [8]. It will completely decay into gravitons (or more generally closed strings), mainly very massive ones, in terms of [7]. The emission of massless closed strings accounts for the Hawking radiation. Part of them can escape from the near-horizon region to infinity. But the massive closed strings can not escape and are left in the black hole. In [8], we have made the estimation that the degenerate states of these emitted closed strings should account for the entropy increased in the host black hole due to absorption of the infalling particle. The result satisfies the first law better for black holes far away from extremality.

Finally, we may need to check the back reaction. For D-particles in string theory, the mass is $m_0 \sim 1/g_s$. Thus, for small $g_s$, the mass of the particle is very heavy. In this case, the particle can deform the vacuum geometry of a black hole strongly on the gravitational theory side while in the tachyon field theory the back reaction is relatively weak [4, 7]. For large $g_s$, the situations are reversed on the two sides. This implies that we may get desired results on one side from controllable calculations on the other side.

# 4 Implications to late-time tachyon cosmology

The above results may have useful implications to tachyon cosmology. The tachyon cosmological model is a powerful model that can provide a unified description of multi-stage puzzles in modern cosmology. However, the reheating mechanism is not well clarified [13]. In what follows, we shall show what clues our above results can provide to this problem and other aspects in this model.

For a flat universe, the FLRW cosmological metric is

$$ds^2 = -dt^2 + a(t)^2 d\vec{x} = a(\eta)^2(-d\eta^2 + d\vec{x}), \tag{23}$$

with $dt/a(t) = d\eta$. The tachyon cosmological model is described by:

$$S = S_{HE} - \int d\eta d\vec{x} a^4 V(T)\sqrt{1 - a^{-2}(\partial_\eta T)^2}. \tag{24}$$

where $V(T) = \tau_3/\cosh(\beta T)$ and $\tau_3$ is the tension of the unstable D3-brane. This action includes the effective theory of a uniform tachyon, which means that the results in previous sections for the Minkowski case are probably applicable here. But the situation is somehow different in this case: the tachyon field is coupled to gravity. This can be tackled by studying the detailed evolutionary process in the model.

From the action, we can have the equations:

$$H^2 = \frac{1}{3M_P^2}\frac{V}{\sqrt{1-\dot{T}^2}}, \tag{25}$$

$$\frac{\ddot{T}}{1-\dot{T}^2} + 3H\dot{T} + (\ln V)' = 0, \tag{26}$$

where the dots and prime respectively denote derivatives with respect to time $t$ and the field $T$.

It is known that the evolution of these equations at early stage leads to an inflationary stage driven by the potential $V$ around $T = 0$. The dS universe also has an event horizon. For a general inflationary universe which is nearly the dS space with $a(t) = e^{Ht}$, there is an equivalent Gibbons-Hawking temperature arising from the periodic Euclideanised time coordinate

$$T_{dS}^{(t)} = \frac{H}{2\pi} \tag{27}$$

This means that the perturbations produced from inflation are nearly Gaussian.

During the inflationary stage, most energy of the potential $V(0)$ is transformed to the rapid expansion of the universe. However, beyond inflation, the situation should be changed. At late stages, the potential quickly evolves into the exponential form: $V(T) \simeq 2\tau_3 e^{-2\beta T}$ and the field solution is simply $T \simeq \text{cosnt} + t$. We can analyze the evolutionary features after inflation by rewriting Eq. (26) at late times as

$$\ddot{T} = (\beta - 3H\dot{T})(1 - \dot{T}^2). \tag{28}$$

For given equation state $\omega \neq -1$, the scale factor evolves as $a(t) \sim t^{2/3(1+\omega)}$. So $H \sim t^{-1}$, i.e., it decreases with time after inflation. Thus, at late times, we still approximately have the solution $\dot{T} = \tanh(\alpha t)$ with $\alpha$ being close to but less than $\beta$, i.e., very close to the solution in flat space. This means that most energy should be transferred to particle creation and radiation at late times. In this case, we can adopt the results obtained in the previous section as a good approximation here. Comparing the action (24) in cosmological background with the standard effective action (49) in Minkowski spacetime, we can read out the temperature for thermal radiation from the tachyon in cosmology

$$T_T^{(\eta)} = \frac{1}{a} T_T^{(t)} \simeq \frac{\beta}{2\pi a}. \tag{29}$$

This temperature cools down as $a(t)$ grows, evolving in the same rule as other thermal components in the expanding universe. It should be the scale of the CMB temperature at the present era. This is like the case for a black hole, the Hawking or Hagedorn temperature observed in the $t$ coordinate is redshifted to a small value as detected for an observer at infinity in the $\chi$ coordinate.

Hence, in the tachyon cosmology, the quantum effects of tachyon can lead to a hot universe after inflation via the particle creation process like what happens near a black hole. In the stringy language, open strings are excited and created from the homogeneous tachyon at the Hagedorn temperature. Meanwhile, closed strings are formed and radiated. The massless closed strings may account for a gravitational background of Gaussian distribution, as can be tested by analysing the tensor perturbations in the next section. The massive closed strings forms the tachyon matter, which is a candidate to dark matter.

This process to attain a hot universe resembles the tachyonic preheating mechanism in hybrid inflation during the stage before the scalar rolls down to the minimum of the potential [34].

We now compare the the temperature for the inflationary relic and the one for radiation from tachyon at late times. As shown in Eq. (44), $\beta$ is dimension of mass: $\beta = M_s/2$ for bosonic strings and $\beta = M_s/\sqrt{2}$ for superstrings. During the slow-roll inflation, we can work out $\dot{H} \simeq -(\beta^2/6)\tanh^2(\beta T)$. As in [13], we can determine that the inflation should end at $H \sim \beta$ from the slow-roll condition $H^2 \gg |\dot{H}|$. Hence, the temperature (27) around this moment is comparable to the Hagedorn temperature on the tachyon

$$T_{dS}^{(t)} \sim T_T^{(t)}. \tag{30}$$

This means that the emission of gravitational waves from the tachyon and the absorption of gravitational waves from the inflationary relic by the tachyon should be balanced in the late-time tachyon cosmology. The situation is also similar to the case in a black hole, where the Hagedorn temperature is equal to the Hawking temperature of the black hole as well.

# 5 Perturbation analysis

In this section, we provide some evidences to arguments made in the previous section via perturbative analysis. We shall simply compare the perturbations in late-time tachyon cosmology with those in a dS universe, since the latter has been well studied and has been proved to deserve thermal features.

It is known that the fields and their fluctuations in dS universes usually satisfy some featured equations, which lead to Gaussian spectra. For example, for a scalar $\phi$ with a potential $V(\phi)$ propagating in the dS universe with the exponential scale factor $a(t) = e^{Ht}$, the perturbative equation is

$$\ddot{\delta}_{\vec{k}} + \left(\vec{k}^2 e^{-2Ht} + \frac{1}{2}V'' - \frac{9}{4}H^2\right)\delta_{\vec{k}} = 0, \tag{31}$$

by separating $\phi(t,\vec{x}) = \phi(t) + \delta\phi(t,\vec{x})$, where $\delta\phi(t,\vec{x}) = a^{-3/2}\sum_{\vec{k}}\delta_{\vec{k}}(t)e^{i\vec{k}\cdot\vec{x}}$. From analysis on the equation, the Gibbons-Hawking temperature (27) can be derived.

In what follows, we show that the perturbations in late-time tachyon cosmology still satisfy dynamical equations similar to (31) (Of course, at early times of tachyon cosmology, we also have the perturbative equations like (31) because the universe is at the tachyon inflationary stage around the top of the potential).

(1) We first consider the tachyon field fluctuations in the cosmological background at late times. We separate the tachyon field as follows: $T(t,\vec{x}) = T(t) + \tau(t,\vec{x})$. In the expanding universe, the perturbed field equation at late times is:

$$\ddot{\tau} + [2\beta\dot{T} + 3H(1-3\dot{T}^2)]\dot{\tau} - (1-\dot{T}^2)a^{-2}\nabla^2\tau = 0. \tag{32}$$

As analysed in the previous section, we approximately have the tachyon solution $\dot{T} \sim \tanh(\beta t)$ at late times, which is also the tachyon solution in Minkowski spacetime. At the end of inflation and afterwards, $H$ will decrease from the value $H \sim \beta$. So we may ignore the second term in the square bracket in the equation. Doing the expansion: $\tau = e^{-\beta t}\sum_{\vec{k}} \tau_{\vec{k}}(t) e^{i\vec{k}\cdot\vec{x}}$, we can get the equation at late times:

$$\ddot{\tau}_{\vec{k}} + \left(\frac{\vec{k}^2}{a^2}e^{-2\beta t} - \beta^2\right)\tau_{\vec{k}} = 0. \tag{33}$$

In the tachyon driven cosmology, $a$ evolves with time. But we can think that it evolves slowly at late times and it can be viewed as a constant within a short time interval. So its behaviour is very like that in Eq. (31).

(2) We now consider the scalar perturbations of the metric in the tachyon cosmology: $ds^2 = -(1+2\Phi)dt^2 + (1-2\Phi)a^2 d\vec{x}^2$. The perturbative equation is given in [12]

$$\varphi_{\vec{k}}'' + \left[(1-\dot{T}^2)\vec{k}^2 - \frac{z''}{z}\right]\varphi_{\vec{k}} = 0, \tag{34}$$

where primes denote derivative with respect to the conformal time $\eta$. The functions are defined as follows

$$z = \frac{\sqrt{3}a\dot{T}}{\sqrt{1-\dot{T}^2}}, \tag{35}$$

$$\frac{\varphi_{\vec{k}}}{z} = \Phi_{\vec{k}} + \frac{2}{3\dot{T}^2}\left(\Phi_{\vec{k}} + \frac{\dot{\Phi}_{\vec{k}}}{H}\right), \tag{36}$$

where $\Phi_{\vec{k}}$ are the coefficients in the Fourier expansion of $\Phi$.

Doing the redefinitions

$$\varphi_{\vec{k}} = a^{-\frac{1}{2}}u_{\vec{k}}, \quad z = a^{-\frac{1}{2}}y, \tag{37}$$

we can express the equation (34) as

$$\ddot{u}_{\vec{k}} + \left[(1-\dot{T}^2)\frac{\vec{k}^2}{a^2} - \frac{\ddot{y}}{y}\right]u_{\vec{k}} = 0. \tag{38}$$

At late times, $y \sim \sqrt{3}a^{3/2}e^{\beta t}$. So $\ddot{y}/y \simeq \beta^2$, which leads to

$$\ddot{u}_{\vec{k}} + \left( \frac{\vec{k}^2}{a^2} e^{-2\beta t} - \beta^2 \right) u_{\vec{k}} = 0. \tag{39}$$

(3) The tensor perturbations accounting for gravitational wave production are described by

$$\ddot{h}_{ij} + 3H\dot{h}_{ij} - \frac{1}{a^2}\vec{\nabla}^2 h_{ij} = 0. \tag{40}$$

Similarly, expand $h_{ij} = a^{-3/2}\sum_{\vec{k}} v_{\vec{k}}(t)e^{i\vec{k}\cdot\vec{x}}\delta_{ij}$. The equation becomes

$$\ddot{v}_{\vec{k}} + \left[ \frac{\vec{k}^2}{a^2} - f(\dot{T}) \right] v_{\vec{k}} = 0, \tag{41}$$

where

$$f(\dot{T}) = \frac{3\ddot{a}}{2a} + \frac{3}{4}H^2 = \frac{3}{4M_P^2} V\sqrt{1-\dot{T}^2}. \tag{42}$$

Since $V(T) \simeq 2\tau_3 e^{-\beta t}$ and $\dot{T} \simeq \tanh(\beta t)$ at late times, we have $f \simeq (3\tau_3/M_P^2)e^{-2\beta t}$.

# 6 Conclusions

In this work, we have shown that the infalling process in a non-extreme black hole background should be a tachyon condensation process. Based on it, we suggest that the Hawking radiation from a black hole could be interpreted as the thermal radiation from rolling tachyon on an infalling probe which behaves as an unstable (D-)particle near the black hole. The Hawking temperature is actually the Hagedorn temperature in the tachyon field theory. We further show that the result is useful in tachyon cosmology. The uniform tachyon rolling in cosmology can automatically create thermal particles to reheat the universe after inflation, via a mechanism just like the Hawking radiation from a black hole. This also means that the perturbations in tachyon cosmology should be still Gaussian even at late times, including the matter-dominated stage. These features may make the tachyon cosmological model still more interesting.

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# Appendix

# A Thermal particle creation from rolling tachyon

Open-string creation from rolling tachyon at the Hagedorn temperature have been discussed in the BCFT. Here, we give a brief introduction to the results and show that similar thermal radiation in the effective field theory are expected.

## A.1 BCFT

In BCFT, the worldsheet action [20, 21] is deformed by the boundary term

$$S_{bndy} = \int_{\partial\Sigma} d\tau m^2(X^0(\tau)), \tag{43}$$

where $m^2(X^0)$ is $(\lambda/2)e^{\pm 2\beta X^0}$ (or their combinations), describing the decay/creation of an unstable D-brane respectively. The constant $\beta$ is

$$\beta = \begin{cases} \frac{1}{2l_s}, & \text{(bosonic string)} \\ \frac{1}{\sqrt{2}l_s}, & \text{(superstring)} \end{cases} \tag{44}$$

with the string length $l_s = \sqrt{\alpha'}$. From the action (43), the non-vanishing components of the energy-momentum tensor at late times are obtained [20, 21]:

$$T_{00} = \frac{1}{2}\tau_p[1+\cos(2\lambda\pi)], \quad T_{ij} \propto \tau_p e^{-2\beta t}\delta_{ij} \tag{45}$$

where $\tau_p$ is the tension of the unstable D$p$-brane and $1 \le i, j \le p$. So we get a pressureless matter state towards the end of condensation, which corresponds to non-relativistic, massive closed strings [22, 23].

The open-string pair creation is discussed in this theory in the minisuperspace approximation [1, 3]. In this approximation, we only consider the zero-mode $x^\mu = (t, \vec{x})$ of $X^\mu = (X^0, \vec{X})$, with the fields $x^\mu$ only depending on $\tau$. Then the action (43) is

$$S = \int d\tau \left[-\frac{1}{4\alpha'}\partial_\tau x_\mu \partial_\tau x^\mu + (N-1) + 2m^2(t)\right], \tag{46}$$

where $N$ is the oscillation number. From this action, we can get the Hamiltonian and further get the Klein-Gordon (KG) equation from the constraint $H = L_0 + \widetilde{L}_0 = 0$:

$$\left\{\partial_t^2 - \vec{\nabla}^2 + \frac{1}{\alpha'}[2m^2(t) + (N-1)]\right\}\psi(t, \vec{x}) = 0, \tag{47}$$

where $m^2(t) = (\lambda/2)e^{2\beta t}$ for brane decay.

The equation (47) can be transformed into the standard Bessel equation. So it can be solved by Bessel functions, providing the ingoing modes, and Hankel functions, providing the outgoing modes. The vacuum with no particles is chosen to be in state in the far past $t \to -\infty$. The production rate of open-string pairs by evaluating the Bogolubov coefficients relating the in and out states are found to be a thermal spectrum at the Hagedorn temperature at late times:

$$T_{\text{Hag}}^{(t)} = \frac{\beta}{2\pi}. \tag{48}$$

The same temperature can also be obtained in the other cases: $m^2(t) = (\lambda/2)e^{-2\beta t}$ and $m^2(t) = \lambda \cosh(2\beta t)$.

This temperature essentially arises from the periodicity of the Euclideanised time coordinate $-it \sim -it + 2\pi/\beta$. It is easy to notice that the equation (47) is invariant under the periodic translation of the Euclideanised time.

The Hagedorn temperature implies that the energy of tachyon will transfer to closed string emission, as discussed in the BCFT in [6, 7]. The results indicate that unstable D-particle or wrapped D$p$-branes will decay into closed strings, mainly very massive ones, towards the end of tachyon condensation. This leads to remarkable speculations for understanding the fate of a particle collapsing into a non-extremal black hole, as will be revealed later. But we will not talk much about this point in this work.

## A.2 Effective field theory

The tachyon effective theory is believed to be equivalent to the BCFT [24, 7, 23], particularly at late times. Thus, it is expected that the same features about string creation and radiation from tachyon obtained in BCFT should be reproduced in this effective theory. In what follows, we show that there should be a similar thermal spectrum of open-string pair creation at Hagedorn temperature in this theory.

The tachyon effective theory [25, 26] for an unstable D$p$-brane in homogeneous case is

$$S(T) = -\int dt d^p x V(T)\sqrt{1 - \dot{T}^2}, \tag{49}$$

where the potential can be [27]

$$V(T) = \frac{\tau_p}{\cosh(\beta T)}. \tag{50}$$

The potential tends to vanish as $T \to \pm\infty$.

At late time as the tachyon grows large, the potential is $V(T) \simeq 2\tau_p e^{-\beta T}$ with $T \simeq t + \text{const}$. If there is no energy loss from the rolling tachyon, the energy density should be constant. So the energy-momentum tensor is derived

$$T_{00} = \frac{V(T)}{\sqrt{1-\dot{T}^2}} = E, \quad T_{ij} = -\frac{V^2}{E}\delta_{ij}, \tag{51}$$

where $E$ is a constant. Thus, the components evolve as those given in Eq. (45) in BCFT. We get a tachyon matter state eventually as well.

The first equation in Eq. (51) says that the speed $|\dot{T}|$ of tachyon is accelerating from 0 to the speed 1 of light, with the acceleration $\ddot{T} \simeq (4\beta\tau_p^2/E^2)e^{-2\beta T}$ at late times. It is known that particles can radiated from accelerating objects due to the Unruh effect. We now consider the open-string creation process in this effective theory.

From the expression of $T_{00}$, the equation of motion can be expressed as

$$\dot{T}^2 = 1 - \frac{1}{E^2}V^2. \tag{52}$$

This equation can be derived from the vanishing of the Hamiltonian[2] as follows

$$H = \frac{1}{2}\left(p_T^2 - 1 + \frac{1}{E^2}V^2\right), \tag{53}$$

with $\dot{T} = \partial H/\partial p_T = p_T$ and $\dot{p}_T = -\partial H/\partial T$.

Similarly, we have the KG equation from $H = 0$:

$$\left(\partial_T^2 + 1 - \frac{\tau_p^2}{E^2}\text{sech}^2(\beta T)\right)\psi(T) = 0. \tag{54}$$

The 1 can be viewed as the "normalised" given by $-\partial_t^2[e^{\pm it}\psi(T)]$. Once the tachyon rolls down off the top of the potential, it can evolve in two parallel and symmetric worlds: $T \geq 0$ and $T \leq 0$. For later purposes, we choose the time direction to be $t \to \infty$ on the positive $T$ side and to be $t \to -\infty$ on the negative side. Since the tachyon field theory (49) can be a theory in Minkowski spacetime, we can define the positive frequency modes as

$$e^{-it}\psi(T), \quad (T \geq 0) \tag{55}$$

$$e^{it}\psi(T). \quad (T \leq 0) \tag{56}$$

In the following discussion, we omit the time-dependent part of the modes.

[2]This is actually not the Hamiltonian density of the tachyon effective theory in the homogeneous case, which is given by $T_{00}$ in Eq. (51).

This equation is the Legendre equation with the transformation $z = \tanh(\beta T)$. Here, we follow the treatments in [28] with the transformations $z = i\sinh(\beta T)$ and $\psi = \sqrt{\cosh(\beta T)}\hat{\psi}$. Then the equation becomes the Legendre equation:

$$\left[(1-z^2)\partial_z^2 - 2z\partial_z + \frac{\lambda^2}{1-z^2} - \left(m^2 + \frac{1}{4}\right)\right]\hat{\psi} = 0, \tag{57}$$

where $\lambda^2 = \tau_p^2/(\beta^2 E^2) - 1/4$ and $m = 1/\beta$. We only consider the case $E \leq 2\tau_p/\beta$ so that $\lambda$ is real, though the discussion of the imaginary $\lambda$ case can also be implemented.

The normalised solution which is smooth across $T = 0$ is given by

$$\psi^{(0)}(T) = e^{-\frac{i\lambda\pi}{2}\mathrm{sgn}(T)} \frac{|\Gamma(\frac{1}{2}+\lambda+im)|}{\sqrt{2\beta}} \sqrt{\cosh(\beta T)} P^{-\lambda}_{-\frac{1}{2}+im}(i\sinh(\beta T)), \tag{58}$$

where $|\Gamma(a+ib)| = \sqrt{\Gamma(a+ib)\Gamma(a-ib)} = \Gamma(a \pm ib)e^{-i\arg\Gamma(a\pm ib)}$. For large $|T|$, the solution is given by the Legendre polynomials of the second kind. The general normalised solutions for positive and negative $T$ are respectively:

$$\psi^{(+)}(T) = h_+^{(+)} + e^{\pi m}h_-^{(+)}, \quad (T > 0) \tag{59}$$

$$\psi^{(-)}(T) = h_-^{(-)} + e^{\pi m}h_+^{(-)}, \quad (T < 0) \tag{60}$$

where

$$h_+^{(\pm)}(T) = \frac{e^{\frac{-\pi m}{2}}}{|\Gamma(\frac{1}{2}-\lambda+im)|}\sqrt{\frac{\cosh(\beta T)}{\beta\sinh(\pi m)}} Q^{-\lambda}_{-\frac{1}{2}+im}(i\sinh(\beta T)). \tag{61}$$

$$h_-^{(\pm)}T) = \frac{e^{\frac{-\pi m}{2}}}{|\Gamma(\frac{1}{2}-\lambda+im)|}\sqrt{\frac{\cosh(\beta T)}{\beta\sinh(\pi m)}} Q^{-\lambda}_{-\frac{1}{2}-im}(i\sinh(\beta T)). \tag{62}$$

The superscripts $(\pm)$ respectively stand for the positive/negative $T$ side and the subscripts $\pm$ respectively stand for the $\pm m$ mode (note that $-1/2 - im = -[1 + (-1/2 + im)]$). The complex conjugates of the modes satisfy

$$\psi^{(0)*}(T) = \psi^{(0)}(-T), \quad \psi^{(\pm)*}(T) = \psi^{(\mp)}(-T), \tag{63}$$

$$h_+^{(+)*}(T) = h_-^{(-)}(-T) = h_-^{(+)}(T), \quad h_-^{(-)*}(T) = h_+^{(+)}(-T) = h_+^{(-)}(T). \tag{64}$$

These give rise to the spatial part of the negative frequency modes.

From the identity $|\Gamma(1-im)|^2 = \pi m/\sinh(\pi m)$, the asymptotical forms are

$$\psi^{(+)} \to \frac{(-)^\lambda e^{-i\frac{\pi}{4}}}{\sqrt{2}}\left(e^{i\theta}e^{-iT} + e^{-i\theta}e^{iT}\right), \quad (T \to \infty) \tag{65}$$

$$\psi^{(-)} \to \frac{(-)^{\lambda} e^{i\frac{\pi}{4}}}{\sqrt{2}} \left( e^{-i\theta} e^{-iT} + e^{i\theta} e^{iT} \right), \quad (T \to -\infty) \tag{66}$$

where $\theta = \arg[\Gamma(1 - im)\Gamma(1/2 - \lambda + im)]$.

The mode expansions are related via

$$h_+^{(+)} = a\psi^{(0)} + b\psi^{(0)*}, \tag{67}$$

$$h_-^{(-)} = a^*\psi^{(0)} + b^*\psi^{(0)*}, \tag{68}$$

with

$$a = -ie^{\pi m} b^* = \frac{-ie^{\frac{i\lambda\pi}{2}} e^{\frac{\pi m}{2}}}{\sqrt{2\sinh(\pi m)}}. \tag{69}$$

From solutions to Eq. (51), it is known that the tachyon field $|T| \to \infty$ as $|t| \to \infty$ so that $|\dot{T}| \to 1$. We first consider the positive $T$ side, on which the time is directed to $+\infty$. As $t \to \infty$, there are two possible solutions. The solution $\dot{T} \simeq 1$ corresponds to the tachyon condensation (brane decay) process, while the one $\dot{T} \simeq -1$ corresponds to the inverse tachyon condensation (brane creation) process (since the tachyon decreases with time growing). So these are two different evolving channels. For the former, the solution $T - t \simeq$ const means that the positive frequency outgoing modes should take the form: $e^{-it}\psi^{(+)} = e^{-it}h_+^{(+)}$. For the latter, the solution reads $T + t \simeq$ const and the mode is $e^{-it}e^{\pi m}h_-^{(+)}$. Similarly, we can get the similar result for $e^{it}\psi^{(-)}$ as the time $t \to -\infty$ on the negative $T$ side. For the tachyon condensation process, $e^{it}\psi^{(-)} = e^{it}e^{\pi m}h_+^{(-)}$.

Here, the tachyon condensation process is concerned. Since the frequency is taken as unit in Eq. (54), from the relation $aa^* - bb^* = 1$, we get the particle numbers at $T \to \infty$ seen in the open string vacuum $|0>$

$$n_1 = |b|^2 = \frac{1}{e^{1/T_H} - 1}, \tag{70}$$

with $T_H = \beta/(2\pi)$ being the Hagedorn temperature given in Eq. (48).

Similarly, the temperature can be viewed as arising from the periodic translation of the imaginary field $-iT \to -iT + 2\pi/\beta$, though this periodicity is not on the time coordinate (but $T \sim t$ at late times). The equation (54) is invariant under this translational symmetry.

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