Probing Axions with Event Horizon Telescope Polarimetric Measurements

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With high spatial resolution, polarimetric imaging of a supermassive black hole, like M87\textsuperscript{★} or Sgr A\textsuperscript{★}, by the Event Horizon Telescope can be used to probe the existence of ultralight bosonic particles, such as axions. Such particles can accumulate around a rotating black hole through superradiance mechanism, forming an axion cloud. When linearly polarized photons are emitted from accretion disk near the horizon, their position angles oscillate due to the birefringent effect when traveling through the axion background. In particular, the supermassive black hole M87\textsuperscript{★} (Sgr A\textsuperscript{★}) can probe axions with masses $O(10^{-20})$ eV ($O(10^{-17})$ eV) and decay constant smaller than $O(10^{16})$ GeV, which is complimentary to black hole spin measurements.

I. INTRODUCTION

The first ever image of the supermassive black hole (SMBH) M87\textsuperscript{★} by the Event Horizon Telescope (EHT) \cite{1,2} leads us to a new era of black hole physics. The high spatial resolution makes the direct visual observation of an SMBH and its surroundings possible. While it offers a new way to study the most extreme objects in our universe predicted by Einstein’s theory of general relativity, we may wonder what else can we learn, especially for fundamental particle physics, from the rich information extracted from EHT under such an extreme environment.

Axion is a hypothetical particle beyond the standard model (SM), which was originally motivated by the solution of the strong CP problem \cite{4,5} in QCD. Beyond the QCD-axion, axion-like particles (ALPs) also generically appear in fundamental theories \cite{6}, and serve as a viable dark matter candidate. There are many search strategies proposed to look for axions, for example, via their conversion into photons \cite{9,10,11,12}, spectral oscillation/distortion of photons \cite{13,14,15}, nuclear magnetic resonance \cite{16,17}, neutron star mergers \cite{18} or various table-top experiments \cite{19,20,21,22,23,24}.

When the Compton wavelength of an axion is at the same order of a rotating black hole size, the axion is expected to develop a large density near the horizon, forming an axion cloud through the superradiance mechanism \cite{25,26,27,28} (for a review see \cite{29}). Such superradiance process can be tested by black hole spin measurements \cite{30,31,32,33,34,35,36,37}. The high spatial resolution makes the direct visual observation of an SMBH and its surroundings possible. While it offers a new way to study the most extreme objects in our universe predicted by Einstein’s theory of general relativity, we may wonder what else can we learn, especially for fundamental particle physics, from the rich information extracted from EHT under such an extreme environment.

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II. PHOTON POLARIZATION FROM BACKGROUND AXION FIELD

The interaction between axion and photon can be written as

\begin{equation}
- \frac{1}{2} g_{\alpha \gamma} a F_{\mu \nu} \tilde{F}^{\mu \nu}.
\end{equation}

This modifies the equation of motion for a photon propagating in an axion background field and it leads to periodic oscillation of a linearly polarized photon’s position angle \cite{38,39,40}. More explicitly, we assume that the variation of the axion field in space and time is much slower than photon’s frequency, i.e., $\mu \ll \omega_{\gamma}$, where $\mu$ is the...
mass of the axion and \( \omega_{\gamma} \) is photon’s frequency. In the later discussion, it becomes clear that the spacetime is approximately flat in the region we are interested. Thus in Lorenz gauge, the plane wave solution of a photon propagating along \( z \)-axis is

\[
A_{\pm}(t, z) = A_{\pm}(t', z') \exp\left[ -i \omega_{\gamma}(t - t') + i \omega(z - z') \right] \\
= ig a \left[ a(t, z) - a(t', z') \right],
\]

where \( \pm \) denotes the two opposite helicity states with

\[
A_0 = A_3 = 0, \quad A_{\pm} = \frac{1}{\sqrt{2}} (A_1 \mp i A_2).
\]

The leading effect of the axion background field comes from the last term of Eq. (2), which results in a rotation of the position angle for a linearly polarized photon,

\[
\Delta \Theta = g a \Delta a(t_{\text{obs}}, x_{\text{obs}}, t_{\text{emit}}, x_{\text{emit}}) \\
= g a \int_{t_{\text{emit}}}^{t_{\text{obs}}} ds \, n^\mu \partial_\mu a \\
= g a \left[ a(t_{\text{obs}}, x_{\text{obs}}) - a(t_{\text{emit}}, x_{\text{emit}}) \right],
\]

Here \( n^\mu \) is the null vector along the path. Note that this only depends on the initial and final axion field values.

In the following discussion, we consider photons emitted near the horizon of a SMBH and observed by a telescope on Earth, e.g. the EHT. Thus we can safely neglect \( a(t_{\text{obs}}, x_{\text{obs}}) \), since the axion field can be very large surrounding the SMBH due to superradiance. Then Eq. (4) becomes

\[
\Delta \Theta \simeq -g a \Delta a(t_{\text{emit}}, x_{\text{emit}}) \\
= -g a_0(x_{\text{emit}}) \cos \left[ \omega t_{\text{emit}} + \beta(x_{\text{emit}}) \right],
\]

Here \( \omega \) is the oscillation frequency of the axion field. It becomes clear in the next section that the axion cloud is approximately non-relativistic, thus the oscillation frequency of the axion background field, and consequently the position angle oscillation frequency of a linearly polarized photon, is determined by axion mass \( \mu \). The amplitude, \( a_0(x_{\text{emit}}) \), is set by the energy density of the axion cloud at the emission point.

### III. SUPERRADIANCE AND BOSENLOVA

Superradiance happens when the following condition is satisfied,

\[
\omega < \omega_{\gamma} \equiv \frac{a_j m}{2 r_+},
\]

where \( \omega \) is the oscillation frequency of axion field, \( m \) is the azimuthal number, \( a_j \) is the dimensionless spin of the black hole, \( r_+ \) is black hole outer horizon radius. Further in Planck unit \( (G_N = c = \hbar = 1) \),

\[
r_\pm = r_g \left( 1 \pm \sqrt{1 - a_j^2} \right) \quad \text{and} \quad r_g = M \quad \text{with} \ M \text{ being the black hole mass.}
\]

The axion field produced through superradiance forms a bound state with the SMBH as a “gravitational atom”.

Define \( \alpha \equiv \mu M \) in Planck unit. In the limit \( \alpha \ll 1 \), the energy levels can be written as

\[
\text{Re}(\omega) \simeq \left( 1 - \frac{a_j^2}{2 n^2} \right) \mu,
\]

where \( n = n + l + 1 \) is the principal quantum number.

In [31], it was shown that the superradiance process is efficient when the Compton wavelength of the axion \( (\lambda_C) \) is comparable to the size of the rotating black hole. The optimized axion mass is a little bit smaller than the one saturating the superradiance condition in Eq. (6).

\[
\frac{r_+}{\lambda_C} \simeq \alpha \simeq 1/2,
\]

The corresponding exponential growth rate of the axion field is related to the imaginary part of \( \omega \).

\[
\tau_{\text{SR}}^{-1} \simeq \text{Im}(\omega).
\]

The simulation shows that \( \text{Im}(\omega)/\mu \) is between \( 10^{-10} \) to \( 10^{-7} \), and it grows with \( a_J \). It is most efficient to populate the highest possible \( L_Z \) state, i.e., \( m = l \). For \( m < l \), \( \tau_{\text{SR}}^{-1} \) is much smaller. Thus in the following discussion, we take \( m = l \).

The superradiance process keeps producing particles until the self-interaction among these bosons becomes important, which leads to non-linear region [34, 35]. In this region, one needs to take into account the whole potential in the effective Lagrangian

\[
S = \int d^4 x \left[ -\frac{1}{2} \partial_\mu a \partial^\mu a - \mu^2 f_a^2 \left( 1 - \cos \frac{a}{f_a} \right) \right],
\]

where the last term is the axion potential. When the expectation value of axion is small, it contains only a mass term. When the amplitude of the axion field in the axion cloud is comparable to \( f_a \), the self-interaction becomes important. In non-relativistic approximation, the axion wave-function takes the form

\[
a = \frac{1}{\sqrt{2 \mu}} (e^{-i \mu t} \psi + e^{i \mu t} \psi^*),
\]

where \( \psi \) is a slowly varying function. Substituting into Eq. (10), we get

\[
S_{\text{NR}} = \int d^3 x \left( i \psi^* \partial_t \psi - \frac{1}{2 \mu} \partial_\mu \psi \partial^\mu \psi^* - \frac{\alpha}{r} \psi^* \psi + \frac{(\psi^* \psi)^2}{16 f_a^2} \right),
\]

where \( \alpha/r \) is the Newtonian gravitational potential of the black hole and the quartic interaction is the leading term for axion self-interaction. The self-interaction becomes important when these two terms are comparable.

\[
\alpha \simeq \frac{\mu a_0^2}{4 f_a^2}.
\]
In [38, 10], simulations are performed to study the non-linear behavior of the axion cloud. After entering non-linear regime, axion cloud either ends as a bosenova explosion or keeps saturating the non-linear region with a steady outflow. In the former case, the self-interaction makes axion cloud collapse and the axion field value decreases by an $O(1)$ factor. The bosenova lasts typically $\tau_{BN} \approx O(100 \sim 1000)\tau_a$. Then the axion cloud starts to build up again until it reaches the non-linear region at a later time. This timescale is much longer than the short period $\tau_a$. In another case, the loss from the gradual scattering towards the far region balances the extraction of the energy from the black hole. The final state does not trigger bosenova explosion, but rather experiences a steady outflow from the axion cloud. Although the axion cloud may end up in two very different possibilities, i.e. bosenova explosion or a continuous outflow, the simulation in [38, 10] indicates that the axion field amplitude in the most dense region, $a_{\text{max}}$, is always around $O(1)$ of $f_a$.

In [39], the bound for axions satisfying superradiance condition of SMBH is shown to exclude only the region for $f_a > 10^{16}$ GeV. This is because when $f_a$ is large enough, before entering non-linear region, enough amount of angular momentum can be extracted from the black hole, and superradiance condition in Eq. (6) no longer holds. Thus, the existence of a black hole with high angular momentum can be used to constrain the parameter space of axion with suitable mass and large $f_a$. However, in the case of $f_a < 10^{16}$ GeV, the later stage, also the most efficient stage, of superradiance is terminated by the axion self-interaction. Further, during bosenova explosion, $O(1)$ fraction of the axion in the axion cloud falls back to black hole. Thus the accumulated change of angular momentum in this case is unlikely to be large enough to become observable. Instead, we here focus on the polarimetric measurements with high spatial resolution, especially with the EHT, which serves a complimentary search with black hole spin measurements.

Using Eq. (6), the maximal change of the position angle can be written as

$$\Delta \Theta_{\text{max}} \simeq -\frac{bc}{2\pi} \cos [\mu_{\text{emit}} + \beta(|x_{\text{emit}}| = r_{\text{max}})]$$

Here $b \equiv a_{\text{max}}/f_a$, which is an $O(1)$ number as discussed above. A constant $c$ is introduced to link $g_{a\gamma}$ and $f_a$

$$g_{a\gamma} = \frac{\alpha_{\text{em}} N}{4\pi f_a} = \frac{c}{2\pi f_a},$$

in which $\alpha_{\text{em}}$ is the fine-structure constant. $N$ is determined by the species of particles contributing to the axion-photon interaction through the triangle diagram. This is model dependent and we treat $c$ as a free parameter.

### IV. AXION FIELDS PROFILE

In this section, we study the axion field spatial profile. We focus on the black hole vicinity, especially the region with the ring feature presented by the EHT, i.e. $r_{\text{ring}} \approx 5.5 r_g$ [9]. We note that the current results published by the EHT group do not have polarization information, but the polarization data is expected to be available in the future [54].

The general solution for scalar field with mass $\mu$ in the Kerr background can be written as

$$a(x^\mu) = e^{-i\omega t} \epsilon^{\mu\phi} S_{lm}(\theta) R_{lm}(r),$$

where $x^\mu = [t, r, \theta, \phi]$ in Boyer-Lindquist coordinates. The $\theta$ dependence is characterized by spherical harmonics

$$S_{lm} = S_l^m \left(\cos \theta, a_j M^2 \sqrt{\omega^2 - \mu^2} \right)$$

This simplifies to spherical harmonics $Y_l^m$ in the non-rotating limit or non-relativistic limit.

Imposing ingoing boundary condition at the outer horizon $r_+$ and setting to zero at infinity, the solution of the radial part can be written as

$$R(r) = (r - r_+)^{-i\sigma} (r - r_-)^i\sigma + \chi^{-1} e^{iqr} \sum_{n=0}^{\infty} a_n \left(\frac{r - r_+}{r - r_-}\right)^n$$

with

$$\sigma = \frac{2r_+ (\omega - \omega_c)}{r_+ - r_-}, \quad q = -\sqrt{\mu^2 - \omega^2}, \quad \chi = \frac{\mu^2 - 2\omega^2}{q}.$$  

Here $r_\pm = r_g \left(1 \pm \sqrt{1 - a_j^2}\right)$ are outer/inner horizon radius. The expansion coefficients, $a_n$, are solved using Leaver’s nomenclature in [31].

This is the solution of a free scalar field and one may worry about the non-linear effect discussed in the last section, which generalizes Klein-Gordon equation to Sine-Gordon equation with non-perturbative potential, as in Eq. (10). This subtlety is studied in the Appendix of [38] where the deviation from Eq. (16) is calculated with the Green’s function method. It is shown that a perturbative transition to other modes that are not satisfying superradiance condition is only significantly induced when there is a bosenova. However, during bosenova, these additional modes fall back to black hole after each collapse and one gets a state similar to the initial perturbative state. Since the superradiance timescale $\tau_{SR}$ is much longer than $\tau_{BN}$, the solution of Klein-Gordon equation is valid in most of time during the large period between each bosenova.

In Fig. 1 we show axion field profile of $l = 1, m = 1$ state. We take $\alpha = 0.4$ and $a_j = 0.99$. At $r_{\text{ring}} = 5.5 r_g$, the axion field value is not significantly different from the maximal value, i.e. $R(r_{\text{ring}}) \approx 0.9 R(r_{\text{max}})$. Taking $b = c = 1$, the position angle variation is between $\pm 8^\circ$. 


FIG. 1: The absolute value and the complex phase of $R(r)$ for $(l = 1, m = 1)$ state. We take $\alpha = 0.4$, $a_J = 0.99$ and $r_g = M$ in Planck unit.

For the time just before the bosenova explosion, it can even reach $\pm 25^\circ$.

Note that the complex phase in $R(r)$ is almost a constant for $r > 2r_+$. Also $S_{lm}$ does not contribute to a large complex phase in the non-relativistic limit. Thus the space-dependent complex phase in Eq. (5) is dominated by $m\phi$ in Eq. (16)

$$\beta(x_{\text{emit}}) \simeq m\phi.$$ (20)

In Fig. 2, assuming that the rotation axis of the disk points to the observer, we show $\Delta \Theta$, at a fixed time, as a function of position. A $\phi$-dependence of the position angle was predicted by the magnetohydrodynamic model of the accretion flow [55]. Therefore, both time dependence of the position angle at a fixed spatial point and the position angle as a function of position at a fixed time can be used to test the existence of axion cloud by the future high-resolution polarimetric imaging by the EHT.

V. DETECTABILITY

As shown in previous sections, taking $c \sim O(1)$, the maximal oscillation amplitude of the position angle is $O(10^\circ)$. This is expected to be well within the capability of the EHT. The previous observations, with a subset of the EHT configuration and an exposure of tens of minutes [56], measures the position angle at precision of $\sim 3^\circ$. It is reasonable to expect a better precision can be achieved with even shorter exposure time for the updated EHT observations. For Sgr A*, the expected oscillation period is $100 \sim 1000$ s (In Table I, we give a summary of parameters for two SMBHs, M87* and Sgr A*). It might be challenging to have a good enough sampling of the observations within one period which typically requires an exposure of e.g. tens of seconds. The situation for M87* is more promising, due to a substantially longer oscillation period (a few days). The upcoming analysis of the polarization data, particularly for M87*, should be able to provide valuable information about the possible axion superradiance around these BHs.

At last, in order to give a realistic estimation, one needs to take into account the spatial resolution of the EHT. The integral along $r$ and $\theta$ directions are always constructive, especially when the accretion disk is concentrated at $\theta = \pi/2$ plane. Only the average on $\phi$, as shown in Eq. (20), may wash out the position angle variation. The spatial resolution of the M87* image is about 20$\mu$as (full width at half maximum; FWHM) [1], which corresponds to a region of $\sim 2r_g$ at a distance of $\sim 17$ Mpc. Assuming a nearly face-on disk of the emission, which is similar to the case of M87* with an inclination angle around $17^\circ$ [57], this spatial resolution translates to $\Delta \phi \simeq 4r_g/r_{\text{ring}} = 0.7$ rad. Without losing generality, taking $\phi = 0$ and considering the average effect within $\Delta \phi$, one obtains the wash-out factor caused by the spatial
resolution as
\[
\int_0^{\Delta \phi} \cos(\mu t + m \phi) d\phi = \frac{\sin(m \Delta \phi/2)}{m \Delta \phi/2} \cos(\mu t + m \Delta \phi/2).
\]  

Before EHT, due to the limitation on spatial resolution, \( \Delta \phi = 2\pi \), and the change of the position angle gets totally averaged out. For EHT, taking \( m = 1 \) and \( \Delta \phi \sim 0.7 \), the wash-out effect is not significant.

### Table I: Typical parameters of the axion superradiance of the two SMBHs, M87* and Sgr A*.

| SMBH   | \( M \)     | \( a_J \) | \( \mu \) range                  | \( \mu \) for \( \alpha = 0.4 \) | \( \tau_a \) | \( \tau_{SR} \) |
|--------|-------------|-----------|----------------------------------|----------------------------------|-------------|---------------|
| M87*   | \( 6.5 \times 10^8 M_\odot \) | 0.99      | \( 2.1 \times (10^{-21} \sim 10^{-20}) \) eV | \( 8.2 \times 10^{-21} \) eV | \( 5.0 \times 10^5 \) s | \( > 1.5 \times 10^{12} \) s |
| Sgr A* | \( 4.3 \times 10^6 M_\odot \) | —         | \( 3.1 \times (10^{-19} \sim 10^{-17}) \) eV | \( 1.2 \times 10^{-17} \) eV | \( 3.3 \times 10^3 \) s | \( > 1.0 \times 10^8 \) s |

Except for the change of position angle, the trajectories of photons with different helicities are also bent differently when they propagate through the axion cloud. This leads to additional subtleties when one measures the change of polarization as a function of time. The bending angle is proportional to \( \mu g_{\gamma}/\omega_\gamma \). In [35], it is shown that the bending angle of a 1 GHz photon is \( O(10^{-4}) \) as the SMBH populates the axion cloud with \( \mu \approx 10^{-18} \) eV. However, the EHT observation uses 230 GHz radio waves, which makes the bending angle smaller than 1\( \mu \)as for both M87* and Sgr A*.  

Thus we conclude that the light bending effect, given the current spatial resolution of EHT, is negligible.

The emission from the accretion disk is also expected to be stable, which makes the identification of the axion-induced signal somehow challenging. Nevertheless, the astrophysical variability of the disk is usually non-periodic. As illustrated in [56], the position angle of the linearly polarized emission shows intra-hour variabilities in the vicinity of a few Schwarzschild radii. However, the variabilities are quite diverse for different observation time. Therefore, the unique behavior of periodic oscillation of the position angle due to the axion field is potentially detectable with high-precision measurements by the EHT, through e.g., a periodicity search after Fourier transforming the data in time domain. On the other hand, [56] shows the Day 82 observation with no clear detection of variabilities. This can be used to constrain the existence of the axion cloud from superradiance.

In Fig. 3, we show the axion parameter space which is potentially probed by M87* and Sgr A*, assuming \( b = 1 \). Notice that this method is complementary to the constraints from black hole spin measurements [35], which potentially excludes the region of large \( f_a \). The switch-over point is determined by the choice of \( c \). A non-observation of periodic oscillation angle will put an upper bound on the value of \( c \) and rule out the corresponding mass window below \( f_a < 10^{16} \) GeV.

### VI. Conclusion and Discussion

Dense axion cloud can be induced by rapidly rotating SMBH through superradiance. The position angle of a

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1 In [35], the light source is assumed to be behind the SMBH while we are considering the photons from the accretion disk, which sits in the axion cloud. Further, the photon bending considered in [35] is along \( \phi \) direction, but we expect the bending remains the same order of magnitude along other directions.
linearly polarized photon emitted near horizon oscillates periodically due to the existence of the axion cloud. The change of the position angle can be tested both temporarily and spatially. A polarimetric measurement with good spatial resolution by e.g., the EHT, is particularly crucial for such a test.

Our proposed search strategy is complimentary to black hole spin measurements where the axion self-interaction cannot be too strong. In addition, since the axion cloud enters the non-linear region, a drastic bosonova or a steady outflow gives \( a_{\text{max}}/f_a \sim O(1) \) for most of the time. Thus our observable does not reply on the detailed dynamics of the axion cloud.

The main model-dependent factor comes from the translation from \( 1/f_a \) to \( g_{a\gamma} \), which is the relation between the scale for non-perturbative potential and that of axion-photon coupling. With \( c \equiv 2\pi f_a g_{a\gamma} \sim O(1) \), the change of the position angle due to axion can be \( O(10^5) \), which is within the capability of current/future EHT measurements.

At last, we note that the position angle oscillation induced by the axion background does not depend on photon frequency. This is a unique property distinct from the Faraday rotation induced by the galactic magnetic field, where the position angle is proportional to the square of photon wavelength. Polarimetric measurements at different frequencies in the future can thus be used to distinguish astrophysical background and improve the sensitivity on testing the axion superradiance scenario.

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