Stringent constraints on axion-photon coupling with Event Horizon Telescope polarimetric measurements of supermassive black hole M87∗

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Introduction

With the technique of very long baseline interferometry, the Event Horizon Telescope (EHT) reaches an unprecedented spatial resolution in the radio frequency band and images the horizon-scale structure of the supermassive black hole (SMBH) M87∗ [1–3]. Such an achievement provides us with unique opportunities to study physics in extreme conditions. Particularly, the EHT measures the electromagnetic radiation at the 1.3 millimeter wavelength, which is optimal for studying the magnetic field around black holes via polarization measurements. Such measurements have been performed very recently [4], and the magnetic field structures near the event horizon were revealed for the first time [5].

Beyond the astrophysics in the vicinity of a SMBH such as the accretion flows and the jet, the polarization measurement can also help us probe new physics beyond the Standard Model in particle physics. In [7], the authors proposed that the temporal and spatial variations of the electric vector position angle (EVPA) around the photon ring can be used to study the existence of an axion cloud produced by the superradiance mechanism, thanks to the axion induced birefringent effect.

More explicitly, the EHT collaboration [5] presented the fiducial polarimetric images for four days (April 5, 6, 10, 11 in 2017). They also provide the azimuthal distribution of the EVPA, including systematic uncertainties, for two days (April 5 and 11). Based on the search strategy proposed in [7] and the information provided in [5], we propose an improved search strategy using differential EVPA in the time domain. Such a method can effectively reduce the uncertainties from the astrophysical background. In addition, we include a more realistic modeling of the accretion flows as well as the radiative transfer calculation of polarized light.

We demonstrate that the newly released EHT polarimetric measurements provide a powerful probe to studying the axion-photon coupling for axions with mass around 10⁻²⁰ eV. Interesting constraints, exceeding the
Black Hole Superradiance} Through the superradiance mechanism, a rapidly spinning black hole can generate an exponentially growing axion cloud, if the axion has a Compton wavelength $\lambda_c$ comparable to the gravitational radius of the Kerr black hole $r_g$ [8,14] (for a review see [15]). Such a process is terminated by either the axion self-interaction or a sufficient loss of the angular momentum of the black hole. The wave-function of the axion cloud can be written as

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

where $x^\mu = [t, r, \theta, \phi]$ are the Boyer-Lindquist (BL) coordinates. Such a bound state formed by the black hole and the axion cloud is a close analogue to a hydrogen atom with coupling constant being $\alpha \equiv r_g/\lambda_c$. The $\theta$ dependence is characterized by the spheroidal harmonics $S_{lm}$ which simplifies to $Y_{lm}$ in the non-rotating limit of the black hole or non-relativistic limit of the axion cloud. $R_{lm}(r)$ is the radial component of the wave function that vanishes at both the horizon of the black hole and the infinity. For a benchmark axion with Compton wavelength satisfying $\alpha = 0.4$, the cloud peaks close to the photon ring, i.e. $r_{\text{ring}} = 5.5 r_g$ [7]. For smaller value of $\alpha$, the radius at which the axion field reaches the maximum, $r_{\text{max}}$, scales as $1/\alpha^2$ and becomes larger.

The highest superradiant rate happens for $l = 1, m = 1$, which is the ground state among those states satisfying the superradiant condition [13]. In this case, among all values of $\theta$, the axion cloud wave function peaks at the equatorial plane of the black hole, i.e. $\sin \theta = 1$. $r_{\text{max}}$ becomes larger for a bigger azimuthal number $l$ with a much longer superradiant timescale. Thus in this study we only focus on $l = m = 1$ state.

The range of superradiant condition for $\alpha$ is sensitive to the black hole spin $a_J$ which is still uncertain for M87*. In Refs. [16,17], M87* was claimed to be a nearly extreme Kerr black hole. On the other hand, $a_J = 0.5$ is more favored than $a_J = 0.94$ in [5] while the latter still satisfies the polarimetric constraints. Thus in this study we take the black hole spin $a_J$ as 0.99 for a benchmark, and discuss its impact in Appendix I.

With the growth of the axion cloud, the axion field value gets close to the decay constant $f_a$ that governs the self-interaction of axions from the axion potential $V(a) = m_a^2 f_a^2 (1 - \cos a/f_a)$, where $m_a$ is the mass of the axion. Due to the nontrivial self-interaction, the superradiance is terminated, and the axion cloud would enter a non-linearly self-interacting regime. The fate of the axion cloud can either be a violent bosenova or a saturating phase [18,22]. Interestingly, in either case, the numerical simulations [18,20] and the analytic estimation [21] show that the maximum of the field value $a_{\text{max}}$ remains close to $f_a$ as long as the saturating regime is reached. This leaves our conclusion insensitive to this subtlety. In the following discussion, we introduce $b \equiv a_{\text{max}}/f_a$ to describe the peak value of the axion cloud.

**Axion induced birefringent effect** Due to the axion-photon coupling $g_{a\gamma}aF_{\mu\nu}F^{\mu\nu}/2$, the temporal and spatial variations of an axion background field induce a change to the dispersion relation of photons. In the Lorentz gauge, the modified Maxwell’s equation for photons propagating along the $z$-direction is

$$\Box A_\pm = \pm 2ig_{a\gamma}[\partial_\tau A_\pm - a\partial_\phi A_\pm],$$

where $\pm$ denotes two opposite helicities with $A_0 = A_3 = 0, A_\pm = (A_1 \mp iA_2)/\sqrt{2}$. Thus the EVPA of a linearly polarized photon, labeled as $\chi$, experiences a rotation due to the axion background as

$$\Delta \chi = g_{a\gamma}[a(t_{\text{obs}}, x_{\text{obs}}) - a(t_{\text{emit}}, x_{\text{emit}})].$$

Note that the EVPA change only depends on the axion field values at the emitting and observing points [23]. A generalization to the curved spacetime was discussed in [24], reaching the same conclusion. We emphasize that such a simple expression only holds when photon propagates in the vacuum and the source size is much smaller than the Compton wavelength. This can be a good approximation when the accretion flows are geometrically thin disk. However, the SMBHs Sgr A* and M87* are both geometrically thick and optically thin at 230 GHz [4,23,26]. Therefore, the medium effect needs to be included by solving the radiative transfer along the photon path. This will be discussed in detail later.

We introduce a parameter $c$ which translates the axion photon coupling $g_{a\gamma}$ to the decay constant $f_a$,

$$g_{a\gamma} \equiv \frac{c}{2\pi f_a},$$

The axion density around the Earth is negligible compared with that of an axion cloud surrounding M87*. Thus for a linearly polarized photon emitted at $(t, r, \theta, \phi)$ in the BL coordinate of the black hole, without the medium effect, its EVPA shift can be written as [7]

$$\Delta \chi(t, r, \theta, \phi) \approx -\frac{bcR_{11}(r) \sin \theta \cos [\omega t - m_a \phi]}{2\pi R_{11}(r_{\text{max}})}.$$  

The ratio $R_{11}(r)/R_{11}(r_{\text{max}})$ depends on the axion mass. Here we replace $S_{lm}(\theta)$ with $\sin \theta$, since the axion field is non-relativistic. As discussed previously, the maximal value of the axion field in the axion cloud can be comparable to $f_a$, thus we have $b \sim O(1)$ in the parameter space we are interested in.
Accretion Flows and Radiative Transfer For low-luminosity active galactic nuclei, such as Sgr A* and M87*, the accretion flows are approximately described as radiatively inefficient accretion flows (RIAFs)\cite{25,27}, which are geometrically thick and optically thin. The thickness of the accretion flows is characterized by a dimensionless quantity $H$\cite{28}. For a magnetically arrested disk (MAD), which nicely fits M87*\cite{6}, $H$ is compressed to be 0.05 by the strong magnetic field in the inner region ($\lesssim 10 r_g$) and becomes $\sim 0.3$ in a farther region\cite{29,32}. In this study, we adopt the analytic RIAF model\cite{28} as a benchmark model. We vary $H$ in order to understand the uncertainties induced by the thickness of the accretion flows. As demonstrated later, a different choice of $H$ does not change our conclusion qualitatively.

The Stokes parameters ($I,Q,U,V$) are generically applied to describe the properties of the macroscopic polarization (see e.g.,\cite{33}), in which $I$ is the total intensity, $Q$ and $U$ characterize the linear polarization, and $V$ describes the circular polarization. As mentioned above, the RIAF model is optically thin, and thus the medium effect of photon propagation is important. In particular, the medium effect leads to mixtures among the four Stokes components. The polarized radiative transfer equation can be written as

$$\frac{d}{ds} \begin{pmatrix} I \\ Q \\ U \\ V \end{pmatrix} = \begin{pmatrix} j_I & \alpha_I & \alpha_Q & \alpha_U & \alpha_V \\ j_Q & \alpha_Q & \alpha_I & \rho_V & \rho_U \\ j_U & \alpha_U & -\rho_V & \alpha_I & \rho_Q \\ j_V & \alpha_V & -\rho_U & -\rho_Q & \alpha_I \end{pmatrix} \begin{pmatrix} I \\ Q \\ U \\ V \end{pmatrix},$$

where $s$ is the proper time, $j_{I,Q,U,V}$ are the polarized emissivities, $\alpha_{I,Q,U,V}$ are the absorption coefficients, and $\rho_{Q,U,V}$ are the Faraday rotation and conversion coefficients.

The axion induced birefringence effect can be properly included in the radiative transfer matrix as

$$\rho_V = \rho_{V}^{\text{FR}} - 2g_{a\gamma} \frac{da}{ds},$$

where the first term is the coefficient of the frequency-dependent Faraday rotation from the plasma, and the second term is from the axion field which is a frequency-independent addition along the line-of-sight. $\rho_V$ characterizes the rotation of the phase in $Q + iU$, which is related to the EVPA through

$$\chi = \frac{1}{2} \arg(Q + iU).$$

To properly model the axion-induced birefringence effect, we modify the polarized radiative transfer equation in general relativistic radiation transfer code IPOLE\cite{34,35}, as indicated in Eq.\cite{7}, with analytic RIAF models in which $H = 0.05$ and 0.3 implemented. For the velocity distribution, we choose the sub-Keplerian flows as an approximation. The magnetic field is taken to be vertical as a benchmark model\cite{6}. In addition, we adjust the normalization of the electron density to be $\sim 10^6$ cm$^{-3}$ compared to the original model so that the total flux density in the image at 230 GHz is about 0.5 Jy\cite{14} and the magnetic field strength is consistent with the EHT estimate\cite{6}. The results of a more comprehensive study are under preparation\cite{36}, where we investigate the effects with various choices of parameters in the RIAF model, such as the direction of the magnetic field (e.g., toroidal or radial) and the velocity distributions. Among these choices, we note that the axion-induced EVPA variations are mostly sensitive to the contribution of lensed photons. Varying other parameters do not affect our qualitative conclusion.

On the sky plane, we use the polar coordinates ($\rho, \varphi$). The origin of the coordinate is placed at the black hole center, and the axis with $\varphi = 0$ is taken to be aligned with the direction of jet projection on this plane.

Now we need to map ($\rho, \varphi$) to the BL coordinates of the SMBH. This can be done through the ray-tracing in IPOLE. For an MAD, the dominant emission comes from the region near the equatorial plane of the black hole\cite{19}. Combined with the fact that the brightest region is around $5.5 r_g$ away from the Earth, i.e. M87* rotates clockwise on the sky plane. In addition, for the region of interest in this study, $\rho$ can be approximately mapped to $r$ in the BL coordinates as well, $\rho \approx r$. There are also photons that propagate around the black hole several times before reaching the Earth, due to the lensing effect\cite{37,40} and can influence the polarization as well\cite{41}. For the total intensity they contribute sub-dominantly with $\sim 10\%$ at most. However, as we will see later, if one uses the intensity weighted EVPA as the observable, their contribution is not negligible and can lead to a noticeable washout on the axion induced birefringence. We note that such a washout effect can be largely calibrated if the detailed EVPA distribution is provided at different $r$.

An example to illustrate the effects is presented in Fig.\cite{1}, where the quivers show the EVPA variation. The length of the quiver is proportional to the intensity of the linear polarization, i.e. $\sqrt{Q^2 + U^2}$, and the direction indicates the polarization direction. The white quiver lines show the values of EVPA without the axion induced birefringence. One oscillation period of the axion cloud is equally divided into 16 segments, and the color of each quiver, from red to purple, represents the time evolution.

**EVPA Variation** Now we follow\cite{5} and calculate the the intensity weighted average EVPA as a function of $\varphi$
on the sky plane,

\[ \langle \chi(\varphi) \rangle \equiv \frac{1}{2} \arg \left( \langle Q \times I \rangle + i \langle U \times I \rangle \right), \tag{9} \]

where the intensity weighted average region is between \( r_{\text{in}} \simeq 3r_g \) and \( r_{\text{out}} \simeq 8r_g \) according to [2].

The axion induced birefringent effect leads to an oscillation of \( \langle \chi(\varphi) \rangle \). For \( l = m = 1 \) state of the axion cloud, such a variation can be generically parametrized as

\[ \Delta \langle \chi(\varphi) \rangle = -A(\varphi) \cos[\omega t + \varphi + \delta(\varphi)]. \tag{10} \]

In Fig. 2, we show results from the IPOLE calculations with \( H = 0.05 \) and \( H = 0.3 \) as two representative models. We study the \( \varphi \)-dependence of the relative phase \( \delta \) and the amplitude \( A \) normalized to \( g_{a\gamma}a_{\text{max}} \equiv bc/2\pi \).

We note that \( \delta(\varphi) \) is well approximated by a sinusoidal function. It fits well as

\[ \delta(\varphi) \approx -\omega R \sin 17^\circ \cos \varphi + \delta_0. \tag{11} \]

where \( \delta_0 \) is the arbitrary initial phase of the axion cloud. \( R \) is approximately the radius of the ring, i.e. \( \sim 5.5 \ r_g \), where the dominant emission comes from. This indicates that \( \delta(\varphi) \) is a consequence of the emission time delay due to the \( 17^\circ \) inclination angle between the M87* spin direction and the sky plane. In addition, such a fit is also a reliable verification on the validity of the simple mapping between \( \varphi \) on the sky plane and \( \phi \) in BL coordinate, as discussed in the previous section.

The behavior of \( A(\varphi) \) is more subtle. There are several aspects necessary to be considered. First, there is an important washout effect along the line-of-sight. The accretion flow is optically thin in our study. Thus photons reach the Earth simultaneously were emitted at different times along the line-of-sight. The axion field oscillates at a frequency \( \omega \sim m_c \). Consequently these photons experience different axion cloud phases at their emission. This results in a washout for \( A(\varphi) \), especially when the radiation size is larger than the Compton wavelength \( \lambda_c \) (see the discussion in the Appendix II). A decrease of \( H \) from 0.3 to 0.05 reduces the washout effect slightly. We choose \( H = 0.3 \) as a conservative benchmark in the rest of the discussion.

The dominant washout effect, which also leads to the variation of \( A(\varphi) \) along the \( \varphi \)-direction, comes from the lensed photons that experience a much longer propagation time than the direct emission from the accretion flow. The axion induced EVPA variations for those photons do not add up coherently. For smaller value of \( \alpha \), the Compton wavelength is longer, and the washout effect is less severe. However, in this case \( A(\varphi)/g_{a\gamma}a_{\text{max}} \) becomes smaller (see the black line in Fig. 2), which is mainly due to the factor \( R_{11}(r)/R_{11}(r_{\text{max}}) \), i.e., the emission ring deviates from the axion cloud peak. Finally the smearing due to the finite angular resolution of the EHT observation also leads to a washout effect along

![FIG. 1: An example from the IPOLE simulation. For axion parameters we take \( \alpha = 0.4 \), and \( g_{a\gamma}a_{\text{max}} = 3\pi/4 \) as a benchmark. An analytic RIAF with \( H = 0.3 \), vertical magnetic field, and sub-Keplerian velocity distribution are assumed. The white quiver is the EVPA without the axion induced birefringent effect at each point, and different colors, ranging from red to purple, represent the EVPA time variation when the axion cloud exists. The plus x axis is taken to be aligned with the direction of jet projection on this plane.](image)

![FIG. 2: The relative phase, \( \delta(\varphi)/2\pi \), and amplitude, \( A(\varphi)/g_{a\gamma}a_{\text{max}} \), of the \( \langle \chi(\varphi) \rangle \) variation are shown. The black hole spin is assumed to be 0.99 and the accretion flow is modeled as the analytic RIAF using IPOLE, with the vertical magnetic field and sub-Keplerian velocity distribution. Different choices of \( H \) and \( \alpha \) are shown in the plot. \( \varphi = 0 \) corresponds to the position angle of the jet projection. The behavior of \( A(\varphi) \) is more subtle. There are several aspects necessary to be considered. First, there is an important washout effect along the line-of-sight. The accretion flow is optically thin in our study. Thus photons reach the Earth simultaneously were emitted at different times along the line-of-sight. The axion field oscillates at a frequency \( \omega \sim m_c \). Consequently these photons experience different axion cloud phases at their emission. This results in a washout for \( A(\varphi) \), especially when the radiation size is larger than the Compton wavelength \( \lambda_c \) (see the discussion in the Appendix II). A decrease of \( H \) from 0.3 to 0.05 reduces the washout effect slightly. We choose \( H = 0.3 \) as a conservative benchmark in the rest of the discussion.](image)
the $\varphi$-direction \cite{7}.

**Analysis Strategy** The EHT collaboration has released the fiducial polarimetric images for 4 days, i.e. April 5, 6, 10 and 11, 2017 \cite{5}. For each of the days, it is a map presenting the EVPA distribution as well as the linear polarization intensity. In addition, for two of the four days, April 5 and 11, the intensity weighted EVPA, $\langle \chi(\varphi) \rangle$, as functions of the azimuthal angle are provided. Along the $\varphi$-axis, $\langle \chi(\varphi) \rangle$ is obtained by taking the average of a fan-like slice with an opening angle of $10^\circ$.

As discussed previously, we are looking for an overall pattern of the EVPA variation across the azimuthal angle of the polarized emission. It is natural to expect that the astrophysical background of two sequential days do not change remarkably \cite{3}. Thus we group the 4-day observations into 2 pairs, (April 5, 6) and (April 10, 11). For each pair, we study the variation of $\langle \chi(\varphi) \rangle$ along the azimuthal direction, i.e. $\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle$ where $t_i$ and $t_j$ represent the two sequential days with interval $t_{\text{int}} \equiv t_j - t_i = 1$ day. To investigate whether there is an evidence of axion induced birefringent effect, we compare the $\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle$, obtained in each pair of days, with the model prediction with the axion field.

To calculate $\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle$, one has to take into account of the the observation time in each day, i.e. $t_{\text{obs}} \approx 6$ hours. Using the parametrization in Eq. (10), we have

$$
\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle = \mathcal{A}(\varphi) \int_{-t_{\text{obs}}/2}^{t_{\text{obs}}/2} \text{d}t \left\{ \cos \omega(t_i + \delta t) + \varphi + \delta(\varphi) \right\}
$$

$$
= \mathcal{A}(\varphi) \left[ \cos \omega(t_i + \delta t) + \varphi + \delta(\varphi) \right] \delta t
$$

$$
= \mathcal{A}(\varphi) \sin \omega(t_i + t_{\text{int}}/2) + \varphi + \delta(\varphi)
$$

(12)

where $\mathcal{A}(\varphi) = 2A(\varphi) \sin [\omega_{\text{int}}/2] \sin [\omega_{\text{obs}}/2]/(\omega_{\text{obs}}/2)$. The factor $\sin [\omega_{\text{int}}/2]/(\omega_{\text{obs}}/2)$ represents a washout effect due to the time average during the observation time of each day. The factor $\sin [\omega_{\text{int}}/2]$, $\omega_{\text{obs}}$ represents how much the axion cloud has changed during the time interval of a day. For the parameter space we are interested in, the axion oscillation period is larger than one day. Thus $\sin [\omega_{\text{int}}/2]$ contributes as a suppression factor. At last, we note that one can absorb $\omega(t_i + t_{\text{int}}/2)$ into the phase of the axion cloud $\delta_0$, which is a nuisance parameter in the analysis.

$\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle$ for the two pairs of sequential days are related. First, the $\mathcal{A}(\varphi)$ characterizes the difference in the EVPA variation amplitude for two sequential days. With a good approximation, $\mathcal{A}(\varphi)$ for the two pairs are the same. Furthermore, the phases of the axion field at the two observational times are related to each other by a relative phase difference as $(\omega \times 5$ days). Such a correlation is used in our analysis.

**Data Characterization** In Fig. (8) of Ref. 5, $\langle \chi(\varphi) \rangle$ are constructed using five different methods. The obtained results do not perfectly agree with each other, whose differences represent the possible systematic uncertainties of the analysis. However, we are mainly interested in the time variation of the EVPA. We assume that the analysis systematics can be largely eliminated for this differential study. As a benchmark, we use the results produced by the polsolve method \cite{42} in our analysis.

The reconstruction of $\langle \chi(\varphi) \rangle$ from the data is subtle. For example, there are nontrivial leakages between the polarization modes in the measurements, characterized by the so-called D-term. These can lead to uncertainties in $\langle \chi(\varphi) \rangle$. For some values of $\varphi$ (particularly when the polarized intensity is low), the reconstruction gives ambiguous results, which appear as bifurcations in the figure. In our analysis, we apply a conservative approach through removing the range of $\varphi$ with reconstruction ambiguities. We choose a bin-size of $10^\circ$ in order to reduce the correlations among different azimuthal angles. The uncertainties in these bins are characterized by the width of the stripe. We only keep the bins whose EVPA distribution can be properly fit by a Gaussian function. As a consequence, the left bins have azimuthal angles ranging between $30^\circ$ to $90^\circ$ and $170^\circ$ to $330^\circ$ for the April 5 data, and $30^\circ$ to $310^\circ$ for the April 11 data \cite{51}. Effectively we have 53 bins in total. For each bin, we fit the stripe width by a Gaussian function to obtain an estimation of the measurement error. The errors roughly range from $\pm 3^\circ$ to $\pm 15^\circ$.

In the fiducial linear polarization images, there are slight differences between the two sequential days. It is not clear whether such differences are induced by the intrinsic variation of the accretion flows or by the change of the baseline coverage \cite{5}. Thus in our analysis, we treat the central value of $\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle$ as zero. Furthermore, we assume that the reconstruction uncertainties for the two sequential days remain the same, but uncorrelated. Thus the error bar of $\langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle$ is simply that of $\langle \chi(\varphi, t_i) \rangle$ multiplied with a factor of $\sqrt{2}$.

For each $\omega \approx m_a$, we calculate the likelihood values for different $g_{\alpha\gamma}a_{\text{max}} \equiv k\epsilon2/\pi$. Throughout the analysis, the axion field phase $\delta_0$ is treated as a nuisance parameter and is marginalized over $[0, 2\pi]$. We focus on the axion mass in the region $\alpha \in [0.10, 0.44]$. For a lower mass, the peak of the axion cloud is too far away from the photon ring and the superradiance timescale is longer than the age of the universe. For a higher mass, the $l = m = 1$ mode no longer satisfies the superradiance condition \cite{14}. In Fig. 3, we present the 95% upper bound on the axion-photon coupling, characterized by the value of $c$. Here we assume that the axion cloud has saturated, i.e. $b = 1$. The bound is mass-dependent, and is weaker for smaller axion masses. This is due to a lower value of $R_{11}(r)/R_{11}(r_{\text{max}})$, as well as a stronger suppression from
2 \sin [\omega \text{int}/2] \text{ in Eq. (12). We also compare our constraints with the bounds from CAST [43] and Supernova 1987A [44] assuming } f_a = 10^{12} \text{ GeV in Fig. (3). A region of previously unexplored parameter space is covered by the EHT observations, with the constraints improved by several times to two orders of magnitude for the axion mass window we consider. Notice that the previous bounds are transformed from } g_{a\gamma} \text{ and a larger value of } f_a \text{ up to } 10^{16} \text{ GeV (for } f_a \text{ above this scale, the axion self-interaction is too weak to stop the superradiance) makes our constraints go beyond further.}

FIG. 3: The 95% upper limits (green) on the axion-photon coupling, characterized by \( \log(c) \), derived using the polarimetric observations of M87* [45]. The bounds from CAST [43] and Supernova 1987A [44], assuming \( f_a = 10^{12} \text{ GeV} \), are shown for comparison.

There should be uncertainties from the time-dependent astrophysical background. The time variation of the accretion flows is not well understood, and consequently it is not included in our simulation discussed above. This leads to the major subtlety when we compare the results from the simulation with the observations. For M87* with \( 6.5 \times 10^9 \) solar mass [14], the typical time scale is 5 days for light propagating near the innermost stable circular orbit (ISCO). In addition, for gas in the accretion flows with sub-Keplerian velocity distribution, the associated time scale is approximately one month at \( r \approx 5.5 r_g \). Thus it is reasonable to expect that the astrophysical features at a large scale, comparable to the size of the accretion flows, remains approximately unchanged at the time scale of one day. The differential EVPA for two sequential days, i.e. \( \langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle \), is an ideal observable with the astrophysical uncertainties being highly suppressed.

To qualitatively estimate such an uncertainty, among the azimuthal angle bins selected in our analysis, we find the common ones shared in April 5 and April 11 data, and calculate the differences between the average values of \( \langle \chi(\varphi) \rangle \). Assuming the accretion flow dynamics leads to an approximately linear variation during the 6-day interval, we divide the difference in each bin by a factor of 6. This provides a rough estimation on the variation induced by the astrophysical background. By comparing it with the error bar derived for \( \langle \chi(\varphi, t_j) \rangle - \langle \chi(\varphi, t_i) \rangle \), we find that the astrophysical background uncertainty is generally smaller than that from the EVPA reconstruction. This justifies the validity of our proposed analysis strategy.

Conclusion and Potential Improvements Polarized imaging of the vicinity of SMBHs offers a very unique probe to search for axions, thanks to the EVPA oscillation of linearly polarized photons due to the axion induced birefringent effect. In this study, we introduced a novel method for data analysis which may significantly reduce the astrophysical background. We applied this method to the four days’ polarization measurements on M87* by the EHT [5]. The EHT observations can rule out a region of the axion mass and axion-photon coupling parameter space which is unexplored by previous experiments.

We emphasize that a further optimized search strategy can be applied with improved measurements and analysis of the polarization emission, by e.g., the upcoming EHT observations with higher cadence or the Next Generation EHT (ngEHT) in the future. First, increasing statistics definitely improves the sensitivity. This includes more available pairs of sequential days for our differential analysis and more stable EVPA reconstruction. It also helps if the results can be presented in shorter time segments within one day. Furthermore, since the axion induced birefringent effect is independent on the photon frequency, the polarization measurements at different frequency bands will be extremely helpful to distinguish the axion induced effect from the Faraday rotation in the plasma. The radial distributions of the EVPA are also valuable. The sensitivity can be improved using the full prediction of the axion profile. Finally, the EVPA observed beyond \( r_{\text{ring}} \) is free from the contamination of lensed photons which dominate the contribution to the washout effect. Removing these lensed photons helps to provide a more universal signal prediction, independent of the accretion flow models [38]. All these improvements are within expectations of either the future EHT data release or the ngEHT with an increased baseline coverage [14].

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For different accretion flow models, such as the Standard Accretion and Normal Evolution (SANE) jet may be the dominant source of emission. However, SANE is not favored to describe the accretion flows of M87*.

Notice that in Ref. 5, 0° corresponds to the Northern pole while the jet projection has a −72° position angle. Thus in their figures, 0° corresponds to ϕ = 72° in our definition.
Appendix I: Impact of Black Hole Spin

For a fixed azimuthal mode $m$ and black hole spin $a_J$, the superradiance condition imposes an upper limit on $\alpha$ [14],

$$\alpha < \frac{a_J m}{2 \left(1 + \sqrt{1 - a_J^2}\right)}.$$  \hfill (A-1)

With $m = 1$, $\alpha$ can be at most 0.5 for an extreme Kerr black hole and 0.25 if $a_J = 0.8$. Once the superradiance condition is satisfied, the axion cloud profile is only slightly influenced by the value of $a_J$ with $\alpha$ being fixed [18].

The accretion flows, on the other hand, can be affected by the black hole spin. Taking sub-Keplerian RIAF as the benchmark model, we compare the axion cloud induced time variations on EVPA with different choices of the black hole spin, $a_J = 0.5$ and $a_J = 0.8$ with the same $\alpha < 0.25$ to satisfy superradiance condition [A-1]. It turns out that the EVPA variations remain qualitatively the same with that of $a_J = 0.99$ case. This is due to the fact that the line of sight washout effect becomes negligible for lower value of $\alpha$ with longer Compton wavelength which is discussed in Appendix II.

Appendix II: Line of Sight Washout Effect

Here we investigate in detail how the line integral along the line of sight leads the washout effect. As discussed previously, the axion induced birefringent effect can be taken into consideration through a modification on the radiative transfer matrix. To gain some intuition, we simplify the problem by ignoring the medium effect, but only keeping the source terms and the axion effect. Then the evolution of the linear polarization can be written as

$$d(Q + i U) = d \frac{ds}{ds} = j_Q + i j_U - i 2 g_{\gamma \gamma} \frac{da}{ds} \left(Q + i U\right).$$  \hfill (A-2)

This can be solved generally as

$$Q(s_f) + i U(s_f) = \int_{s_i}^{s_f} e^{i 2 g_{\gamma \gamma} \left(a(s_f) - a(s)\right)} \left(j_Q(s) + i j_U(s)\right) ds,$$  \hfill (A-3)

where $s_i$ and $s_f$ are the initial and final points along the line of sight. For further simplification, we assume the radiation source terms, $j_Q$ and $j_U$, as constants in a finite region. Thus the washout effect is characterized by the axion-dependent integral in Eq. (A-3). For a qualitative estimation, we take $a(s)$ as a coherently oscillating background with a constant amplitude $a_{\text{max}}$ in the same region. Then the washout effect on $A$ defined in Eq. (10) are shown in Fig. (A-1), as a function of the size of the radiation source $S_r$, normalized by the axion Compton wavelength $\lambda_c$. Here we see that the washout effect is negligible if $\lambda_c \gg S_r$, and it becomes sizable when these two scales are comparable.

In a realistic scenario like the optically thin accretion flow around M87*, $S_r$ is determined by the geometric thickness of the accretion flow. The washout effect is not significant if the accretion flow satisfies $r_{\text{ring}} H < \lambda_c$. Taking $\alpha \equiv r_g/\lambda_c = 0.4$ and $r_{\text{ring}} \simeq 5.5 r_g$, this condition becomes $H < \frac{1}{2}$, which is satisfied in our RIAF model. For smaller $\alpha$, due to the much larger $\lambda_c$, this washout effect along the line of sight is further reduced.