The X-Ray Polarization of the Accretion Disk Corona of Active Galactic Nuclei

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Abstract

Hard X-rays observed in Active Galactic Nuclei (AGNs) are thought to originate from the Comptonization of the optical/UV accretion disk photons in a hot corona. Polarization studies of these photons can help to constrain the corona geometry and the plasma properties. We have developed a ray-tracing code that simulates the Comptonization of accretion disk photons in coronae of arbitrary shapes, and use it here to study the polarization of the X-ray emission from wedge and spherical coronae. We study the predicted polarization signatures for the fully relativistic and various approximate treatments of the elemental Compton scattering processes. We furthermore use the code to evaluate the impact of nonthermal electrons and cyclo-synchrotron photons on the polarization properties. Finally, we model the NuSTAR observations of the Seyfert I galaxy Mrk 335 and predict the associated polarization signal. Our studies show that X-ray polarimetry missions such as NASA’s Imaging X-ray Polarimetry Explorer and the X-ray Imaging Polarimetry Explorer proposed to ESA will provide valuable new information about the physical properties of the plasma close to the event horizon of AGN black holes.

Key words: accretion, accretion disks – black hole physics – galaxies: Seyfert – polarization – scattering

1. Introduction

Active galactic nuclei (AGNs) are powerful sources of X-rays. Their spectrum is dominated by a power-law continuum presumably emitted by hot and possibly partially nonthermal plasma of particles, known as a corona. Repeated inverse Compton processes in the corona energize optical/UV photons originating from an accretion disk emitting in the IR/optical/UV. Even though the first accretion disk and corona models were developed in the seventies (e.g., Novikov & Thorne 1973; Shakura & Sunyaev 1973) and refined over the last 40 years (e.g., Haardt & Maraschi 1991; Dove et al. 1997; Nowak et al. 2002), the geometry of the corona, i.e., its location and spatial extent, is still a matter of intense debate (Gilfanov & Merloni 2014).

Recent X-ray reverberation observations (e.g., Fabian et al. 2009; Wilkins & Fabian 2013) and future X-ray polarization observations offer a new way of constraining the corona geometry that is complementary to the more traditional constraints from X-ray spectroscopy. Polarimetric observations with the Imaging X-ray Polarimetry Explorer (IXPE) Small Explorer mission (Weisskopf et al. 2014), and possibly with the proposed X-ray Imaging Polarimetry Explorer (XIPE) ESA mission (Soffitta et al. 2013) promise to provide geometrical information about the corona and the inner structure of accretion disks. The polarization of X-ray emission from the accretion disk of a stellar mass black hole (BH) is predicted to be linear polarized with the polarization fraction being a function of inclination of the disk (Li et al. 2009, and references there in). In the case of AGNs, X-rays are emitting from a hot corona in the vicinity of the accretion disk. The polarization of this emission depends on the scattering off the accretion disk and on the scattering processes in the corona. This dependency of the polarization on the scattering makes polarization studies a promising way to distinguish between different corona geometries. Schnittman & Krolik (2010) showed that in stellar mass BHs, the corona geometry has a major impact on the predicted energy spectra of the polarization fraction and the polarization angle. Dovciak et al. (2012) studied the polarization of unpolarized corona X-rays scattering off the accretion disk of AGNs. In this paper, we study for the first time the impact of the Klein–Nishina (K–N, Klein & Nishina 1929) cross section on the polarization of the coronal emission. Furthermore, we study how nonthermal electrons in the coronal plasma and polarized synchrotron and cyclotron seed photons affect the observable polarization properties. Our studies are based on a general relativistic ray-tracing code that simulates the individual scattering processes accounting for the energy dependent K–N cross section in the framework of a general relativistic ray tracing code. The code assumes that the 3D corona plasma orbits the black hole with the angular velocity of a Zero Angular Momentum Observer (ZAMO). The seed photons are polarized with an initial polarization given by the classical results of Chandrasekhar (1960). The code tracks photons forward in time, making it possible to study repeated scatterings in the corona and off the accretion disk.

Although simple models assume a single-temperature corona, the coronal plasma may have a distribution of temperatures and/or an admixture of nonthermal plasma (e.g., from magnetic reconnection in the corona). The Atacama Large Millimeter/Submillimeter Array may be able to reveal the presence of nonthermal plasma in AGN coronae (Inoue & Doi 2014). The energy spectra of black holes in X-ray binaries show clear evidence for nonthermal particles (e.g., Coppi & Blandford 1990; Poutanen & Coppi 1998; Gierliński et al. 1999; Romero et al. 2014; Malzac 2016). In Cyg X-1, the nonthermal component was detected to be strongly polarized (Laurent et al. 2011; Jourdain et al. 2012), and this was taken as evidence that it is formed as synchrotron emission in the jet rather than as inverse Compton emission in the corona. Below, we will use our code to evaluate the possibility that nonthermal electrons in the corona produce high polarization fractions at high energies.
The structure and strength of the magnetic fields in AGN accretion disks is still a matter of debate (e.g., Blandford & Znajek 1977; Belmont & Tagger 2005). Cyclotron or synchrotron (cyclo-synchrotron) photons are naturally expected from the energetic electrons radiating in the ambient coronal magnetic field (see, e.g., Malzac & Belmont 2009; Veledina et al. 2011). In this paper, we will evaluate if X-ray polarization could contribute to clarifying the situation if a fraction of the seed photons are highly polarized cyclotron or synchrotron photons.

The rest of the paper is structured as follows. We describe the ray-tracing code and the corona geometries in Section 2. We report on the results of the studies of the impact of the K–N cross section, nonthermal electrons, and cyclotron and synchrotron seed photons in Sections 3–5. In Section 6, we model the NuSTAR observations of Mrk 335, and use the modeling to predict the polarization fraction and angle energy spectra. We summarize our results in Section 7. Throughout this paper, all distances are given in units of the gravitational radius \( r_g = GM/c^2 \), and we set \( G = c = \hbar = 1 \). The mass of AGN is \( 10^8 M_\odot \) unless otherwise specified. The inclination is \( i = 0^\circ \) for an observer viewing the disk face-on and \( i = 90^\circ \) for an observer viewing the disk edge-on.

2. Ray Tracing
2.1. Thermal Disk Simulation

We expanded upon the general relativistic ray-tracing code described in Krawczynski (2012), Beheshtipour et al. (2016), Hoormann et al. (2016), called the thermal code in the following. The code simulates an accretion disk extending from the innermost stable circular orbit (ISCO) \( r_{\text{ISCO}} \) to \( r_{\text{max}} = 100 r_g \). The disk emits thermally with a radial brightness distribution given by mass, energy, and angular momentum conservation (Page & Thorne 1974). Our simulations ignore the emission and scattering by gas inside the ISCO. The code uses Boyer–Lindquist coordinates and tracks photons forward in time using the fourth-order Runge–Kutta method to integrate the geodesic equation and to parallel transport the polarization vector. All photons are tracked until they come to within 0.02% \( r_g \) to the event horizon at which point we assume they will cross the event horizon and will not contribute to the observed signal. The code assigns an initial polarization to each photon using Chandrasekhar’s parameterization of the limb brightening and polarization of the emission from an indefinitely deep atmosphere of free electrons (Chandrasekhar 1960, Tables 24 and 25). The polarization fraction is Lorentz invariant and only changes when the photon scatters.

When a photon hits the disk, its wave and polarization four vectors are transformed from the global Boyer–Lindquist coordinate system into the plasma rest frame, the polarization fraction and the polarization vector are used to calculate the Stokes parameters, a random new direction is drawn, Chandrasekhar’s results for Thomson scattering off an indefinitely deep electron atmosphere are used to modify the Stokes parameters for the given incident and scattered direction, the Stokes \( I \) parameter is used to modify the statistical weight of the photon, the \( I, Q, \) and \( U \) parameters are used to calculate the polarization fraction and vector, and the wave and polarization four vectors are back-transformed from the plasma rest frame into the global Boyer–Lindquist coordinate system.

Figure 1. Meridian plane coordinate system used to derive polarization and Stokes parameters.

Chandrasekhar (1960) derived the following equation for the intensities \( I(\tau, \mu) \) and \( I(\tau, \mu) \) in the directions parallel and perpendicular to the meridian plane (see Figure 1) and the Stokes parameter \( U \) (Chandrasekhar 1960, Chapter 10, Equation (163)):

\[
I(0, \mu, \phi) = \begin{pmatrix} I_l \\ I_r \\ U \end{pmatrix} = \frac{1}{4\mu} F \begin{pmatrix} Q(S, \mu, \phi) \\ 0 \\ 1 \end{pmatrix}
\]

where \( \mu_0 \) and \( \mu \) are the direction cosines of the incident and scattered emission, \( F \) and \( I \) are the incident fluxes in the \( l \) and \( r \) directions, and \( Q \) and \( S \) are matrices tabulated in Chandrasekhar (1960). For each scattering, we multiply the statistical weight with \( 2\pi \eta \mu/\mu_0 I/F \) (with \( I = I_l + I_r \)). The factor \( 2\pi \) normalizes the expression to \( \eta \) when averaging over all scattering directions with \( \eta \leq 1 \) being the scattering efficiency (\( \eta = 1 \) if not mentioned otherwise), and the factor \( \mu/\mu_0 \) normalizes \( I \) and \( F \) to the flux per accretion disk area. Note that the polarization fraction is given by \( \Pi = |(I_l - I_r)/(I_l + I_r)| \).

We track photons until they reach a Boyer–Lindquist radial coordinate of \( r = 10,000 r_g \). The photon wave and polarization vectors are subsequently transformed into the reference frame of an observer at fixed coordinates. We study inclination dependent effects by collecting photons arriving within \( \pm 4^\circ \) of the \( \theta \)-angle of the observer. We calculate the polarization fraction \( \Pi \) and angle \( \chi \) by summing the Stokes parameters of the individual photons and then using the standard equations

\[
\Pi = \sqrt{(U^2 + Q^2)/I^2} \quad (2)
\]

\[
\chi = 1/2 \tan^{-1}(U/Q), \quad (3)
\]

with the summed Stokes parameters \( I, Q, \) and \( U \). We find that it is important to calculate statistical errors on the polarization results. We do so by calculating the Stokes parameters for 11 independent subsets of the simulated data sets, and deriving 11 independent estimates of \( I, Q, \) and \( U \). The mean values and standard deviations of the three Stokes parameters are subsequently used to calculate the mean values and standard
deviations of the polarization fraction and angle based on Equations (2) and (3) and standard error propagation.

2.2. Comptonization in the Corona

We simulate isothermal wedge and spherical coronae. The wedge coronae lie above and below the accretion disk and are chosen to have a constant opening angle of $8^\circ$ (see Figure 2(a)). A photon traversing a distance $dl$ in the corona rest frame scatters with the probability $p = e^{-\sigma r}$ with

$$d\tau(r, z) = \kappa \rho(r, z) dl,$$

where $\kappa = 0.4 \text{ cm}^2 \text{ g}^{-1}$ is the opacity to electron scattering and $\rho(r, z)$ is the density of the coronal plasma at radius $r$ and height $z$. For the wedge corona, we assume the coronal gas density dependence on radius and height from (Schnittman & Krolik 2010, Equation (6))

$$\rho(r, z) = \rho_0(r) \exp(-z/H(r)), \quad \rho_0(r) = \frac{\tau_0}{\kappa H(r)}$$

with the adjustable parameter $\tau_0$. Since the total optical depth is a function of the worldline, we characterize the density with the scattering coefficient (the optical depth per proper length) $\sigma = d\tau/dl$, and we quote $\sigma$ for the mean height and the mean radius of the disk.

The spherical corona extends from $r_{SCO}$ to $R_{edge}$. In this geometry, the accretion disk is truncated at the outer edge of the corona at $R_{edge} = 15R_g$ (Figure 2(b)). The optical depth is chosen to be a linear function of radius, $\tau(r) = (\tau_0/R_e)r$ and $R_e = R_{edge} = r_{SCO}$ is the radius of the corona and $\tau_0 = \tau_0/\kappa R_e$. In the wedge corona, all seed photons are thermally emitted accretion disk photons. The disk of the spherical corona model is truncated at the edge of the corona, and we assume the same seed photon luminosity and energy spectrum as function of the radial coordinate $r$ as for the wedge corona, and launch the photons with a random polar angle $\theta$ with a flat $\cos(\theta)$ distribution.

The calculation assumes that the coronal plasma rotates with the ZAMO at an angular velocity of $\Omega_0 = -\Omega_{pol}/\Omega_{ISCO}$ with $\Omega_{pol}$ and $\Omega_{ISCO}$ being components of the Kerr metric. A random number is drawn to decide if the photon scatters. If it does, its wave vector and polarization vector are transformed first to the rest frame of the coronal plasma, and then to the electron rest frame, where the new scattering energy and polarization vector are calculated. This process will be described in more detail in the next section. After the scattering, the photon wave vector and polarization vector are transformed back, first into the coronal rest frame (CR) and then into the global Boyer–Lindquist frame, where the tracking continues.

3. Thomson and K–N Scattering

We simulate photon–electron scatterings using the Thomson approximation and the full K–N cross section. In both cases, we transform the wave vector and polarization vector of the photon first from the global BL coordinates into the corona frame coordinates, and subsequently into the rest frame of one of the scattering electrons (assumed to be isotropic in the corona frame). We randomly draw the direction of the scattered photon in the electron rest frame. In the Thomson approximation, the scattering does not change the photon energy. More accurately, the photon looses energy according to the Compton’s equation,

$$\epsilon_1 = \frac{\epsilon_0}{1 + x(1 - \cos \theta)},$$

where $\epsilon_1$ and $\epsilon_0$ are the energy of the photon after and before scattering, respectively, $x$ is the photon energy in units of the electron rest mass, and $\theta$ is the scattering angle.

We use the Stokes parameters and the nonrelativistic Raleigh and relativistic Fano scattering matrices to calculate the statistical weight of the scattering and polarization fraction of the scattered photon. The Stokes parameters are calculated with the help of two sets of basis vectors (see Figure 3). The projection of the polarization vector onto the first set of basis vectors allows us to calculate the polarization angle $\chi_0$, and the Stokes parameters $Q_0 = \Pi_0 \cos 2\chi_0$ and $U_0 = \Pi_0 \sin 2\chi_0$ with $\Pi_0$ being the polarization fraction of the incoming photon.
The Stokes parameters before and after scattering (subscripts 0 and 1, respectively) are related via

\[
\begin{pmatrix}
    l_1 \\
    Q_1 \\
    U_1
\end{pmatrix} = T_{R/K-N} \begin{pmatrix}
    l_0 \\
    Q_0 \\
    U_0
\end{pmatrix}.
\] (7)

The Raleigh scattering matrix is given by (Chandrasekhar 1960)

\[
T_R = \frac{1}{2} \rho_0^2 \begin{pmatrix}
    1 + \cos^2 \theta & \sin^2 \theta & 0 \\
    \sin^2 \theta & 1 + \cos^2 \theta & 0 \\
    0 & 0 & 2 - \cos \theta
\end{pmatrix}
\] (8)

and the expression for the Fano scattering matrix reads (Fano 1957; McMaster 1961)

\[
T_{K-N} = \frac{1}{2} \rho_0^2 \begin{pmatrix}
    \epsilon_1 \\
    \epsilon_0
\end{pmatrix}^2 \begin{pmatrix}
    1 + \cos^2 \theta + \frac{1}{m_e c^2} (\epsilon_0 - \epsilon_1) (1 - \cos \theta) & \sin^2 \theta & 0 \\
    \sin^2 \theta & 1 + \cos^2 \theta & 0 \\
    0 & 0 & 2 - \cos \theta
\end{pmatrix},
\] (9)

The Stokes parameters give use the polarization fraction $\Pi_1$ and angle $\chi_1$ after scattering according to Equations (2) and (3), and we use $\chi_1$ to calculate the polarization vector $f'$ of the outgoing photon. We transform the wave vector and polarization vector back into the corona frame and the global BL coordinates.

For each scattering, we multiply the statistical weight of the photon with a factor $w_1 = l_1/l_0$ encoding the physics of the scattering process in the rest frame of the scattering electron, and with the kinematical factor $w_2 = (1 - \beta \cos(\theta_0))$ with $\theta_0$ being the angle between the photon’s and electron’s momentum vectors in the coronal rest frame. The factor $w_2$ reflects the higher likelihood of photon–electron head-on collisions compared to photon–electron tail-on collisions. Because the factor is frequently omitted in simulations of Compton interactions, we briefly justify it based on the derivation of the rate of Compton scatterings of electrons immersed in an isotropic bath of photons (Blumenthal & Gould 1970; Zdziarski et al. 1991; Bicknell 2017). In the rest frame of a scattering electron, the number of scatterings is given by

\[
\frac{dN'}{dt'} = c_{\sigma T} \int f'(p') d^3p',
\] (10)

with $c$ being the speed of light, $c_{\sigma T}$ the Thomson cross section, $p'$ the momentum of the scattered photons, and $f'(p')$ the number density of photons per momentum volume element $d^3p'$. The scattering rate can be transformed into the rest frame of the coronal plasma (undashed variables) noting that $dN$ and $f(p)$ are Lorentz scalars, $dt' = 1/\gamma dt$, and $d^3p' = (1 - \cos(\theta))d^3p$ with $\beta$ being the electron’s velocity in units of the speed of light in the corona rest frame. Assuming isotropic photons in the rest frame of the coronal plasma $f(p) = f(p)$, the scattering rate is

\[
\frac{dN}{dt} = c_{\sigma T} \int (1 - \beta \cos(\theta))f(p) d^3p,
\] (11)

demonstrating that scatterings with pitch angles $\theta$ contribute with a weight proportional to $1 - \beta \cos(\theta)$.

Figure 4 shows the energy spectra obtained for simulating a spherical corona with $\tau_0 = 3$ with four different treatments: (i) in the Thomson approximation (the photon energy does not change in the rest frame of the scattering electrons), using the Raleigh scattering matrix and omitting the weighting factor $1 - \beta \cos(\theta)$; (ii) same as (i) but including the weighting factor $1 - \beta \cos(\theta)$; (iii) same as (ii) but accounting for the energy loss of the photon in the rest frame of the scattering electrons, and (iv) same as (iii) but using the Fano scattering matrix. We see that the weighting factor makes a noticeable difference and changes the 1–10 keV photon index by $\Delta \Gamma \approx 0.1$. Replacing the Raleigh scattering matrix by the Fano scattering matrix does not noticeably impact the energy spectrum in 1–10 keV; while at higher energies it changes the photon index by $\Delta \Gamma \approx 0.4$.

Figure 5 compares the polarization fraction energy spectra for treatments (iii) and (iv) and shows that, using the proper K–N cross section instead of the Thomson cross section, neither impacts the polarization properties of the keV photons. Interestingly, the K–N cross section shows a difference when increasing the energy of the seed photons by a factor of 50 (adequate for the accretion disks of accreting stellar mass black holes), see Figure 6. In these plots, corona densities (scattering coefficients) are set to give the same spectral index for both treatments (iii) and (iv) in 2–10 keV. For such high seed photon energies, photons scatter fewer times to get into the X-ray band, increasing the importance of each individual scattering. The polarization differences are larger for the wedge corona compared to the spherical corona because, for the former, scatterings are generally more important, as the corona covers a larger fraction of the inner accretion disk area. Based on Equation 6, we expect more pronounced K–N effects at the
highest energies. However, the results reveal significant differences at $<10$ keV energies. As the K–N cross section leads to a reduced scattering rate at higher energies, more photons end up in the $<10$ keV band. The higher polarization fraction in K–N scatterings explains the higher polarization of the $<10$ keV photons.

Figure 7 shows the polarization fraction for the same spectral index of the two geometries for different scattering coefficients. We choose corona densities (scattering coefficients) giving the same net 2–10 keV spectral index of the Comptonized emission for the two geometries. This reveals that at $\Gamma > 1.1$, lower $\tau_0/\sigma$, the wedge corona is more polarized than the spherical corona; while at $\Gamma = 1.1$, the two geometries almost have the same polarization fraction; and at a steeper spectrum the spherical corona is more polarized at certain bins. This result shows that the difference between the geometries becomes even somewhat more significant for smaller $\Gamma$-values. The polarization angle of the two models are almost the same at higher energies, while at lower energies for all spectra the two geometries show opposite polarization directions, Figure 7(b).

In the wedge corona, photons are influenced by two types of scattering, scattering in the corona and scattering off the disk, while in the spherical corona photons that scatter off the accretion disk are much less than in the wedge corona, so the scattering is dominant by the corona scattering. The angle of polarization for the photons scattered off the disk is in the opposite direction of the coronal scattering (Schnittman & Krolik 2010). At lower energies, the two coronae show opposite polarization angles because of the difference in the dominant scattering in the two geometries. While at higher energies mostly coronal scattering is dominated in the two geometries, the polarization angles of the two models are almost the same. In practice, it is hard to distinguish models based on the polarization direction alone, because the orientation of the spin axis of the accretion disk is not well constrained observationally.

As mentioned above, the simulations assume that the wedge corona orbits the black hole with the angular frequency of a ZAMO. As a consequence, the photons originating in the disk will experience Compton scatterings owing to the bulk motion of the coronal plasma relative to the accretion disk. We studied the effect of this bulk motion by running additional simulations with a corona corotating with the underlying disk (called Keplerian corona in the following). Figure 8 compares the polarization fraction of the ZAMO and Keplerian coronas—all other parameters being equal. The results show that ZAMO coronas give higher polarization fractions than Keplerian coronas, owing to the bulk motion of the former. At higher energies, the effect is not noticeable because the large number of scatterings reduce the impact of the first few scatterings on the net polarization signal.

It is important to note that the two coronae cover different portions of the accretion disk with different thermal photon properties. To assess the impact of this point, we performed the same analysis using only seed photons coming from the same portion of the accretion disk for both models. The results were the same as in Figure 7. In our result, the increase in the polarization fraction when increasing the optical depth/scattering coefficient is not as clear as Schnittman & Krolik (2010), Figures 6 and 15. They increased $\tau$ and adjusted the...
corona temperature to maintain the same Compton y parameter. Thus, they compare different corona scattering coefficients for the same energy flux. However, in our case, the temperature of the corona is constant and we are comparing polarization for different spectral energies.

4. Nonthermal Plasma

In this section, we assess the impact of a nonthermal power-law component of the coronal electron plasma on the observed polarization properties. The energy spectra of AGNs and stellar mass black holes in X-ray binaries often exhibit evidence for the presence of such a component (see Johnson et al. 1997; Coppi 2003; Malzac & Belmont 2009; Inoue & Doi 2014, and references therein). We assume that the energy spectrum of the thermal electron component is described by the Maxwell–Juttner distribution:

\[
\frac{dN}{d\gamma} = \frac{\gamma^2 \beta}{\theta K_5 \left( \frac{1}{\theta} \right)} \exp \left( -\frac{\gamma}{\theta} \right),
\]

where \( \gamma \) is the Lorentz factor of the electrons, \( \theta = kT_e/m_e c^2 \) is the electron temperature in units of the electron rest mass, and \( \beta = v_e/c \) is the electron velocity in units of the speed of light. The Maxwell–Juttner distribution describes plasmas with temperatures exceeding 100 keV. At lower temperatures, the distribution resembles the nonrelativistic Maxwell distribution. The nonthermal electron component is given by \( dN/d\gamma = \xi \gamma^{-p} \) with the normalization constant \( \xi \) and the power-law index \( p \). (Oriented Scintillation Spectrometer Experiment) observations of NGC 4151, Johnson et al. (1997) estimated...
that \( \sim 8\% \) but less than 15\% of the source power is in the nonthermal electron component. The results agree with those of Fabian et al. (2017) who studied several NuSTAR AGN observations and estimate that the nonthermal component carries 10\%–30\% of the source power. In the following, we assume a nonthermal energy population with \( \gamma \)-factors between 1 and 1000 carrying 15\% of the total energy of the coronal electrons.

Figure 9 shows the impact of the nonthermal component on the observed polarization signatures. We assume \( p = 3 \) and a rather low temperature of the thermal component of 50 keV so that we can see the impact of the nonthermal electrons at the high-energy end of the observed energy spectra. For both corona geometries, the addition of a nonthermal electron component carrying 15\% of the energy barely changes the polarization properties. Figure 10 shows additional details for the spherical corona model. Photons scattering only off nonthermal electrons are highly polarized as a small number of scatterings results in a high polarization fraction. Photons scattering only off thermal, and off thermal and nonthermal electrons exhibit very similar polarization properties.

Figure 11 shows the polarization of photons as a function of their arrival direction. Figure 11(a) shows photons only scattering off nonthermal photons with a high polarization fraction (encoded by the length of the black bars) when they originate from the inner part of the accretion flow \( (r < 15r_g) \). Figure 11(b) shows the much smaller polarization of photons scattering only off thermal electrons, and Figure 11(c) shows that the polarization of all photons very much resembles the results of Figure 11(b). The polarization angle for all images of photons coming from the corona is very comparable and the patterns are very similar. However, the polarization fractions are higher for image (a) explaining the larger bars. Outside of the corona, photons scatter off the disk, so they are more vertically polarized. Inside the corona, the spherical symmetry of the corona and the axial symmetry of the background metric and accretion disk lead to a spherical polarization pattern.

Ghisellini et al. (1993) argue that the coronal plasma may not have sufficient time to thermalize. Figure 12 compares a thermal model with \( T_e = 171 \) keV with a nonthermal model with \( p = -2 \) for \( 1 < \gamma < 3 \). We choose the same corona
temperature and power-law index as in Ghisellini et al. (1993), giving the same spectral index in 2–10 keV. The two models produce almost identical polarization energy spectra with the difference being most pronounced at photon energies exceeding 500 keV.

### 5. Cyclo-synchrotron Seed Photons

Depending on the magnetic field strength, the electrons may loose a good fraction of their energy by emitting cyclo-synchrotron photons. In this section, we explore the impact on the X-ray polarization energy spectra. We assume an ordered magnetic field of strength $B$ oriented either perpendicular to the disk along the z-axis, or parallel to the disk in the x and y plane. The cyclo-synchrotron photons are partially circularly and partially linearly polarized. Since X-ray polarimeters can only measure the linear polarization, we neglect the circular polarization in the following. As above, we assume the presence of a power-law electron component with $dN/d\gamma = \xi \gamma^{-p}$ for electrons emitting cyclo-synchrotron photons. The synchrotron photons are polarized perpendicular to the magnetic field with a polarization fraction of (Rybicki & Lightman 1979)

$$\Pi_0 = (p + 1)/(p + 7/3). \quad (13)$$

Nonrelativistic electrons emit cyclotron photons with a polarization fraction of

$$\Pi_0 = \frac{1 - \cos^2(\theta)}{1 + \cos^2(\theta)}, \quad (14)$$

where $\theta$ is the angle between the magnetic field and the line of sight (Rybicki & Lightman 1979). In the following, we consider a scenario in which 50% of the seed photons are thermal photons from the accretion disk (polarized according to Chandrasekhar’s equation), and 50% are synchrotron photons from an electron power-law distribution with $p = 3$ for which Equation (13) gives a polarization fraction of 75%.

Figure 13 shows the polarization energy spectra for the spherical corona geometry and the magnetic field parallel to the accretion disk. The red line shows the polarization fraction of the synchrotron photons. As the synchrotron photons are emitted in the infrared band and only a tiny fraction makes it into the X-ray band via a large number of Compton scatterings, the error bars on the polarization fraction and direction are rather large. Accordingly, the X-ray emission is strongly dominated by the thermal seed photons even if a substantial fraction of the seed photon energy goes into the synchrotron component. We obtain the same results for the magnetic field being perpendicular to the accretion disk. Accounting for the effect of synchrotron self-absorption reduces the overall impact of the synchrotron seed photons on the observed energy spectra.
even more. As cyclotron photons are less energetic and exhibit lower polarization fractions than synchrotron photons, we expect that they have a similarly negligible impact on the observed polarization energy spectra. Our result of a negligible impact of the synchrotron seed photons on the emitted energy spectra agrees with the earlier findings of Schnittman & Krolik (2013).

6. Simulations of Seyfert I Galaxy Mrk 335

In this section, we present the results of our code when used to model the *NuSTAR* observations of the Seyfert I galaxy Mrk 335 Keek & Ballantyne (2016). The source harbors a supermassive BH with a mass of $M = 2.6 \times 10^7 M_\odot$ accreting at a rate of $\dot{M} = 0.2 \dot{M}_{\text{edd}}$ (Wilkins et al. 2015). Fitting the *NuSTAR* energy spectrum gives a $2\sim10\text{ keV}$ photon index of $\Gamma \approx 1.9$, and the fit of the Fe Kα line suggests a black hole spin of $a = 0.89$ in geometric units and a black hole inclination of $\sim 70^\circ$. For AGNs, the accretion disk inclination can be inferred from fitting the Fe K-alpha line, or from a combined fit of the flux and polarization energy spectra.

For each corona geometry, we choose three different corona sizes, and adjust the optical depth to recover the observed spectral index. Figure 14 shows the observed and simulated
energy spectra, Figure 15 shows the polarization fraction, and Table 1 lists the model parameters. The shaded area in the plots shows the energy band that IXPE and XIPE can measure in the near future. In this energy range, there is a clear difference between the two corona geometries, while it is hard to constrain the size. The average polarization fraction of Mrk 335 in the energy of 2–10 keV, the IXPE energy band, as a function of inclination is shown in Figure 16. In the wedge corona, one can clearly see the polarization difference between the inclinations, while it is hard to compare in the spherical. The difference between the two geometries is also clear in different inclinations. Mrk 335 is a faint source with the flux of about 10^{-11} \text{ erg cm}^{-2} \text{s}^{-1}, which will need two to three days of IXPE observation for an MDP of 10%. For brighter sources, e.g., NGC 4151 with one order of magnitude higher fluxes (e.g., Marin et al. 2016), IXPE will achieve an MDP of 10% in less than two hours and an MDP of 3% in one day.

7. Discussion

The results presented above can be summarized as follows. The simulated X-ray polarization energy spectra depend strongly on the proper treatment of the kinematic effects (change of the photon energy in the rest frame of the scattering electrons, and relative probability for head-on and tail-on collisions) and the use of the relativistic cross section. Using
the K–N cross section rather than the Thomson cross section increases the 1–10 keV polarization fractions by as much as ∼3% for the wedge corona of the hot accretion disk of a stellar mass BH. For the colder accretion disks of AGNs, the cross section does not impact the predicted polarization properties noticeably. The difference between the different corona geometries depends on the optical depth of the coronal plasma, or, conversely, on the energy spectrum of the observed emission. For high optical depths and hard energy spectra, the spherical corona emission is more polarized than the wedge corona emission in the 2–20 keV energy band. For small optical depths and soft energy spectra, the wedge corona emission exhibits a 1%–2% higher polarization than the spherical corona over the entire 1–100 keV energy range. The different polarization fractions are accompanied by differences in the polarization direction.

We find that a nonthermal electron component with about 15% of the internal electron energy has a negligible impact on the observed polarization properties. Even a completely nonthermal corona with a small optical depth producing the same energy spectrum as a thermal corona shows similar polarization properties as a fully thermalized corona. Similarly, cyclotron and synchrotron photons do not impact the polarization energy spectra strongly—even when they carry a substantial fraction of the seed photons’ luminosity. The reason is that these cyclo-synchrotron photons are expected to have too long wavelengths so that only a small fraction is scattered into the X-ray energy range. Synchrotron photons could have an impact if generated with much shorter wavelengths, e.g., as a consequence of magnetic reconnection events in the corona.

Finally, we presented simulations for the Seyfert I galaxy Mrk 335. We predicted the polarization for the two corona geometries and three different corona sizes. Keeping the energy spectrum fixed, the different corona models predict different polarization fractions and angles. We anticipate that the upcoming IXPE mission will add valuable observables for constraining the properties of the inner engine of AGNs—in particular, if several complimentary techniques including spectral, timing, and polarization analyses can all be used for one and the same object.

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