On High-energy Particles in Accretion Disk Coronae of Supermassive Black Holes: Implications for MeV Gamma-rays and High-energy Neutrinos from AGN Cores

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Abstract

Recent observations with ALMA have revealed evidence for nonthermal synchrotron emission from the core regions of two nearby Seyfert galaxies. This suggests that the coronae of accretion disks in active galactic nuclei (AGNs) can be conducive to the acceleration of nonthermal electrons, in addition to the hot, thermal electrons responsible for their X-ray emission through thermal Comptonization. Here, we investigate the mechanism of such particle acceleration, based on observationally inferred parameters for AGN disk coronae. One possibility to account for the observed nonthermal electrons is diffusive shock acceleration, as long as the gyrofactor $\eta_g$ does not exceed $\sim 10^8$. These nonthermal electrons can generate gamma-rays via inverse Compton scattering of disk photons, which can appear in the MeV band, while those with energies above $\sim 100$ MeV would be attenuated via internal $\gamma\gamma$ pair production. The integrated emission from all AGNs with thermal and nonthermal Comptonization can reproduce the observed cosmic background radiation in X-rays as well as gamma-rays up to $\sim 10$ MeV. Furthermore, if protons are accelerated in the same conditions as electrons and $\eta_g \sim 30$, our observationally motivated model is also able to account for the diffuse neutrino flux at energies below 100–300 TeV. The next generation of MeV gamma-ray and neutrino facilities can test these expectations by searching for signals from bright, nearby Seyfert galaxies such as NGC 4151 and IC 4329A.

Key words: acceleration of particles – accretion, accretion disks – black hole physics – galaxies: active – quasars: supermassive black holes

1. Introduction

Active galactic nuclei (AGNs) are powered by mass accretion onto supermassive black holes (SMBHs). They emit intense electromagnetic radiation in a broad range of frequencies. Measurements of X-ray spectra of AGNs allow us to study various aspect of SMBHs, such as black hole spins (e.g., Reynolds 2014), geometrical structures (e.g., Ramos Almeida & Ricci 2017), and cosmological evolution (e.g., Ueda et al. 2014).

A key for understanding these phenomena is primary X-ray radiation of the accretion disk, which arises from Comptonization of disk photons in moderately thick thermal plasma, namely coronae, above an accretion disk (e.g., Katz 1976; Bisnovatyi-Kogan & Blinnikov 1977; Pozdnyakov et al. 1977; Galeev et al. 1979; Takahara 1979; Sunyaev & Titarchuk 1980). X-ray observations have indicated a coronal temperature of $\sim 10^9$ K and a Thomson scattering opacity of $\gtrsim 1$ (e.g., Zdziarski et al. 1994; Fabian et al. 2015). However, the nature of AGN coronae is still veiled in mystery.

Very recently, Inoue & Doi (2018) reported the detection of coronal radio synchrotron emission from two nearby Seyferts (e.g., Di Matteo et al. 1997; Inoue & Doi 2014; Raginski & Laor 2016), utilizing the Atacama Large Millimeter/submillimeter Array (ALMA). The inferred coronal magnetic field strength was $\sim 10$ G, with a size of $40 R_\odot$, where $R_\odot$ is the Schwarzschild radius, for both active SMBHs with a mass of $\sim 10^8 M_\odot$. It was also found that coronae of Seyferts contain both thermal and nonthermal electrons. This implies that acceleration of high-energy particles occurs in AGN coronae.

High-energy particles in the nuclei of Seyferts have long been discussed in the literature. In the past, it was argued that primary X-ray emission comes from pair cascades induced by high-energy particles accelerated in and/or around accretion flow (e.g., Kazanas & Ellison 1986; Zdziarski 1986; Ghisellini et al. 2004). In the pair cascade model, particles are accelerated by shock dissipation in accretion flows (e.g., Cowie & Lee 1982; Protheroe & Kazanas 1983; Kazanas & Ellison 1986; Zdziarski 1986; Sikora et al. 1987; Begelman et al. 1990). However, the detection of the AGN spectral cutoffs (e.g., Madejski et al. 1995; Zdziarski et al. 2000) and nondetection of Seyfert AGNs in the gamma-ray band (e.g., Lin et al. 1993) ruled out the pair cascade scenario as a dominant source for the primary X-ray emission.

In this paper, we investigate the production mechanism of the observed high-energy particles in AGN coronae. As an example, we consider those high-energy particles supplied by diffusive shock acceleration (DSA) processes (e.g., Drury 1983; Blandford & Eichler 1987) in the coronae. In contrast to the previously discussed AGN accretion shock models, a much lower shock power is required to explain the observed nonthermal species and remain in concordance with the current picture of coronal X-ray emission. Moreover, previous studies

6 High-energy particles in the corona of X-ray binaries have been also discussed in the literature (e.g., Bhattacharyya et al. 2003, 2006).
7 TeV gamma-rays are measured from the Galactic center (HESS Collaboration et al. 2016). This detection indicated possible particle acceleration in accretion flow, even though accretion rate in the Galactic center is several orders of magnitude lower than that in standard disks.
of high-energy particles in AGN accretion disks have treated corona size and magnetic field, which are important parameters for the understanding of particle acceleration, as free parameters. The ALMA observations have allowed us to determine both of them (Inoue & Doi 2018). Most critically, the observationally determined strength of the magnetic field appears to be significantly smaller than the one previously considered in the literature. We take into account these newly determined coronal parameters.

Thermal coronal emission from Seyferts is known to explain the entirety of cosmic X-ray background radiation (e.g., Ueda et al. 2014). In contrast, the origin of the cosmic MeV gamma-ray background radiation from 0.1 MeV to several tens MeV is still unknown (e.g., Inoue 2014). Here, the nonthermal electrons in coronae seen by ALMA will invoke power-law MeV gamma-ray emission via Comptonization of disk photons. Such nonthermal emission has been suggested as a possible explanation for the cosmic MeV gamma-ray background radiation (Inoue et al. 2008). However, nonthermal electron species in the previous work were included in an ad hoc way. In this work, we revisit the contribution of Seyferts to the MeV gamma-ray background radiation by considering the particle acceleration of nonthermal populations in coronae together with the latest X-ray luminosity function of Seyferts (Ueda et al. 2014).

High-energy particles around accretion disks of AGNs also generate intense neutrino emission through hadronuclear (pp) and photomeson ($p\gamma$) interaction processes via interactions between accreting gas and photon fields (e.g., Eichler 1979; Begelman et al. 1990; Stecker et al. 1992; Alvarez-Muñiz & Mészáros 2004). Although these originally predicted fluxes have been significantly constrained by high-energy neutrino observations (IceCube Collaboration 2005), recent studies have revisited the estimated fluxes and found that AGN core models are still viable (Stecker 2005, 2013; Kalashev et al. 2015). However, normalization of neutrino fluxes from AGNs and acceleration properties of high-energy particles in those models are assumed to match with observations. In this work, we also discuss the possible contribution from AGN cores, given our ALMA observations, and investigate the parameter spaces required for the explanation of the IceCube diffuse neutrino fluxes.

We describe general particle acceleration processes in AGN coronae in Section 2. The broadband emission spectrum of the central region of AGNs and physical properties of AGN coronae are presented in Section 3. Relevant timescales and steady-state particle spectra are discussed in Sections 4 and 5, respectively. Sections 6 and 7 present the results of the expected gamma-ray and neutrino fluxes from individual AGN cores and the cosmic gamma-ray and neutrino background fluxes from AGN cores, respectively. Discussion, including other possible particle acceleration mechanism is given in Section 8, and conclusions are in Section 9. Throughout this paper, we adopt the standard cosmological parameters of $(h, \Omega_M, \Omega_{\Lambda}) = (0.7, 0.3, 0.7)$.

2. Particle Acceleration in Nuclei of Seyferts

As nonthermal coronal synchrotron emission is seen in nearby Seyferts (Inoue & Doi 2018), particle acceleration should occur in AGN coronae, even though thermal populations are energetically dominant. Particle acceleration mechanism in the coronae is highly uncertain. Various acceleration mechanisms can take place in the coronae, such as the DSA mechanism (e.g., Drury 1983; Blandford & Eichler 1987), turbulent acceleration (e.g., Zhﺫ PKKů et al. 2018), magnetosphere acceleration (e.g., Beskin et al. 1992; Levinson 2000), and magnetic reconnection (e.g., Hoshino & Lyubarsky 2012). In this work, for simplicity, we consider the DSA as the fiducial particle acceleration process. We discuss the other possible acceleration processes in Section 8.3.

In order to investigate particle acceleration mechanism of the observed nonthermal electrons, we consider the interaction of locally injected relativistic particles with the matter, photons, and magnetic field in the infalling coronae. Although the location of shock sites is uncertain, for simplicity, we assume that shocks occur inside of the coronae. The shock accelerates a part of inflow plasma to high energies. As the energy loss timescale of high-energy protons is, in general, longer than the freefall timescale, a sufficiently high energy density of relativistic particles is maintained to provide pressure to support a standing shock around an SMBH (Protheroe & Kazanas 1983).

Coronae are assumed to be spherical, with a radius of $R_c \equiv r_c R_s$. Here, $r_c$ is the dimensionless parameter of the corona size and $R_s = 2G M_{\text{BH}}/c^2$, where $G$ is the gravitational constant, $M_{\text{BH}}$ is the mass of the central SMBH, and $c$ is the speed of light. Coronae are also set to be in a steady state. We also do not consider positrons in coronae. Thus, the proton number density $n_p$ is equal to the electron density $n_e$ in this work, which gives the maximum number of protons in coronae. In this case, $n_e$ is defined through the Thomson scattering opacity in coronae, $\tau_T$ as

$$n_e = \frac{\tau_T}{\sigma_T R_c} \approx 1.4 \times 10^9 \left(\frac{\tau_T}{1.1}\right) \left(\frac{r_c}{40}\right)^{-1} \left(\frac{M_{\text{BH}}}{10^8 M_{\odot}}\right)^{-1} \text{cm}^{-3},$$

where $\sigma_T$ is the Thomson scattering cross section.

2.1. Dynamical Timescale

The gas is assumed to be spherically accreted onto the SMBH with freefall velocity $v_{\text{ff}} = \sqrt{2G M_{\text{BH}}/R_s}$. The freefall timescale from the coronal region is estimated to be

$$t_{\text{fall}} = R_c/v_{\text{ff}} \approx 2.5 \times 10^5 \left(\frac{r_c}{40}\right)^{1/2} \left(\frac{M_{\text{BH}}}{10^8 M_{\odot}}\right) \text{s}.$$  \hspace{1cm} (2)

2.2. Radiative Cooling

High-energy particles lose their energies through radiative cooling processes. In AGN coronae, high-energy electrons mainly lose their energies via synchrotron and inverse Compton (IC) radiation. The synchrotron cooling rate for an electron with a Lorentz factor of $\gamma_e$ is

$$t_{\text{syn},e}(\gamma_e) = \frac{3}{4} \frac{m_e c}{\sigma_T U_B} \gamma_e^{-1},$$

$$\approx 7.7 \times 10^4 \left(\frac{B}{10 \, \text{G}}\right)^{-2} \left(\frac{\gamma_e}{100}\right)^{-1} \text{s},$$

where $m_e$ is the electron rest mass and $U_B = B^2/8\pi$ is the magnetic field energy density of magnetic field strength $B$. 


The inverse Compton cooling rate, including the Klein–Nishina (KN) cross section (Jones 1968; Moderski et al. 2005; Khangulyan et al. 2014), is

$$t_{IC}(\gamma_e) = \frac{3m_e c}{4\sigma_T} \left[ \int_0^\infty \frac{df_{KN}(b)}{b} \frac{U_{ph}(\epsilon)}{\epsilon} \right]^{-1} \gamma_e^{-1},$$

where $b = 4\gamma_e^2/m_e c^2$ and $f_{KN} \approx 1/(1 + b)$ (Moderski et al. 2005). Here, $\epsilon$ is the target photon energy and $U_{ph}$ is the photon energy density given as $U_{ph}(\epsilon) = L_{ph}(\epsilon)/4\pi R_*^2 c$. The total AGN disk luminosity, $L_{ph}$, which includes contribution from the accretion disk and corona, is defined in Section 3.1. For simplicity, we consider a uniform photon density in the coronae. If the coronae has spatially homogeneous emissivity rather than uniform emission, the mean photon density inside the source is enhanced by a factor of $\sim 2.24$, on average (Atoyan & Aharonian 1996).

For the typical characteristics of the coronae, the energy density of the photon field is

$$U_{ph,\text{tot}} = \int d\epsilon U_{ph}(\epsilon) \sim 5 \times 10^3 \frac{L_{ph,\text{bol}}}{2 \times 10^{45} \text{ erg s}^{-1}} \left( \frac{r_g}{40} \right)^{-2} \times \left( \frac{M_{BH}}{10^6 M_\odot} \right)^{-2} (\text{erg cm}^{-3}).$$

For the magnetic field strength inferred with ALMA, $B \approx 10$ G for $M_{BH} = 10^8 M_\odot$ SMBHs, the energy density of the photon field exceeds the magnetic field energy density if $L_{ph,\text{bol}} \gg 2 \times 10^{42} \text{ erg s}^{-1}$. We note that the dominance of photon fields over magnetic field does not necessarily prevent particle acceleration, as such conditions are met in some efficient nonthermal sources, e.g., in gamma-ray binary systems (Aharonian et al. 2006; Khangulyan et al. 2008). Moreover, high density of target photons can enable the converter acceleration mechanism if a relativistic velocity jump is present in the system (Derishev et al. 2003).

Relativistic protons are predominately cooled though inelastic $pp$ interactions, $p\gamma$ reactions, and proton IC/synchrotron channels. Because the Thomson regime is the only one relevant to proton IC cooling, the proton synchrotron and IC cooling timescales are

$$t_{IC/syn,p} = \frac{3}{4} \frac{m_p}{m_e} \frac{m_e c^2}{\epsilon \gamma_p} U_{ph/B} \gamma_p^{-1},$$

where $m_p$ is the proton rest mass and $\gamma_p$ is the proton Lorentz factor. In the case of the synchrotron losses, this yields

$$t_{syn,p} \approx 4.8 \times 10^{14} \left( \frac{B}{10 \text{ G}} \right)^{-2} \left( \frac{\gamma_p}{100} \right)^{-1} (\text{s}).$$

Given the higher energy density of the photon field, the IC cooling time can be up to $\sim 10^4$ times faster. These electrodynamic cooling channels are inefficient, compared to the hadronic mechanisms below. Hereafter, we do not consider proton IC/synchrotron coolings.

The $pp$ cooling time can be expressed as

$$t_{pp} = \frac{1}{n_p \sigma_{pp} \kappa_{pp}},$$

$$\approx 1.6 \times 10^6 \left( \frac{\gamma_p}{11} \right)^{-1} \left( \frac{r_g}{40} \right) \left( \frac{M_{BH}}{10^6 M_\odot} \right) (\text{s}),$$

where $\kappa_{pp} \approx 0.5$ is the proton inelasticity of the process and we adopt $\sigma_{pp} = 3 \times 10^{-26}$ cm$^2$. Below, we adopt the formalism developed by Kelner et al. (2006). The total cross section of the inelastic $pp$ process $\sigma_{pp}$ is represented as a function of the proton energy $E_p = \gamma_p m_e c^2$.

$$\sigma_{pp} \approx (34.3 + 1.88L + 0.25L^2) \left[ 1 - \left( \frac{E_{pp,\text{thr}}}{E_p} \right)^4 \right] \text{mb} \tag{9}$$

for $E_p \geq E_{pp,\text{thr}}$, where $1 \text{ mb} = 10^{-27}$ cm$^2$, $L = \log(E_p/1 \text{ TeV})$, and $E_{pp,\text{thr}} = 1.22 \text{ GeV}$ (Kelner et al. 2006).

The $p\gamma$ cooling time via photomeson interactions is

$$t_{p\gamma}^{-1} = \frac{c}{2\gamma_p^2} \int_{\tilde{\varepsilon}_{\text{thr}}}^{\infty} d\varepsilon \sigma_{p\gamma}(\varepsilon) \kappa_{p\gamma}(\varepsilon) \int_{\varepsilon/2\gamma_p}^{\infty} d\epsilon U_{ph}(\epsilon) \epsilon^{-4},$$

where $\tilde{\varepsilon}$ and $\epsilon$ are the photon energy in the proton rest frame and the black hole frame, respectively, $U_{ph}$ is the energy density of the photon target, and $\tilde{\varepsilon}_{\text{thr}} = 145 \text{ MeV}$. For numerical calculation, we follow the formalism suggested by Kelner & Aharonian (2008).

The $p\gamma$ interaction also generates pairs, via the so-called Bethe–Heitler pair production process, and its cooling timescale is approximated as (Gao et al. 2012)

$$t_{\text{BH}}^{-1} \approx \frac{7(m_p c^2)^3 \alpha_f c}{9\sqrt{2} \pi m_p^2 c^2} \left( \frac{B}{10 \text{ G}} \right) \int_{m_e c^2/\gamma_p}^{\infty} d\epsilon U_{ph}(\epsilon) \epsilon^{-4} \times \left\{ \left( \frac{2\gamma_p c}{m_e c^2} \right)^{3/2} \left[ \log \left( \frac{2\gamma_p c}{m_e c^2} \right)^{-2/3} + 2/3 \right] \right\},$$

where $\alpha_f$ is the fine-structure constant.

2.3. Acceleration

In the framework of DSA (e.g., Drury 1983; Blandford & Eichler 1987), the acceleration timescale can be approximated as

$$t_{\text{DSA}} \approx \frac{\eta_{\text{acc}} D(E_{\text{CR}})}{v_{sh}^2},$$

where $D$ is the diffusion coefficient, $E_{\text{CR}}$ is the particle energy, and $v_{sh}$ is the shock speed. Here, $\eta_{\text{acc}}$ is a numerical factor that depends on the shock compression ratio and the spatial dependence of $D$ (Drury 1983). We set $\eta_{\text{acc}} = 10$. Assuming a Bohm-like diffusion,

$$D(E_{\text{CR}}) \approx \frac{\eta_{k} e E_{\text{CR}}}{3eB},$$

where $e$ is the electric charge and $\eta_{k}$ is the gyrofactor—which is the mean free path of a particle in units of the gyroradius. We use $\eta_{k}$ to characterize the efficiency of the acceleration. A value of $\eta_{k} = 1$ corresponds to the Bohm limit case. The DSA time
can be written as
\[
I_{\text{DSA}} \simeq \frac{10}{3} \frac{\eta_{\gamma} cR_g}{v_{\text{sh}}} \left( \frac{m_p}{m_e} \right) \left( \frac{r_g}{40} \right) \left( \frac{B}{10 \text{ G}} \right)^{-1} \left( \frac{\gamma}{100} \right) (s),
\]
where \( R_g \) is the gyro radius and \( v_{\text{sh}} \) is set as \( v_{\text{th}}(R_g) \). The value of \( \eta_{\gamma} \) varies in different astrophysical environments. For example, \( \eta_{\gamma} \sim 1 \) can be seen in a Galactic supernova remnant (Uchiyama et al. 2007), while \( \eta_{\gamma} \sim 10^4 \) is seen in the case of blazars in the framework of one-zone leptonic models (e.g., Inoue & Takahara 1996; Finke et al. 2008; Inoue & Tanaka 2016).

3. Properties of Active SMBHs

In this section, we summarize the general observational properties of the central region of AGNs related to high-energy particles in coronae.

3.1. Broadband Emission from the Core Region

Emission from the AGN core region mainly arises from two components (Elvis et al. 1994). The first component consists of the geometrically thin and optically thick standard accretion disks (Shakura & Sunyaev 1973). A standard accretion disk generates a big blue bump from optical to UV, attributed to multicolor blackbody radiation. The second component consists of the Comptonized accretion disk photons from the coronal regions above the accretion disk (Katz 1976; Bisnovaty-Kogan & Blinnikov 1977; Pozdniakov et al. 1977; Sunyaev & Titarchuk 1980). This Comptonized emission appears in the X-ray band, along with emission reprocessed by the surrounding cold materials; a so-called Compton reflection component (e.g., Lightman & White 1988; Magdziarz & Zdziarski 1995; Ricci et al. 2011).

In this work, for the primary X-ray emission from coronae, we assume a cutoff power-law model in the form of \( E^{-\Gamma} \exp(E/E_c) \), where we set \( \Gamma = 1.9 \) and \( E_c = 300 \text{ keV} \) (Ueda et al. 2003, 2014). For the Compton reflection component, we use the \texttt{pexrav} model (Magdziarz & Zdziarski 1995), assuming a solid angle of \( 2\pi \), an inclination angle of \( \cos i = 0.5 \), and the solar abundance for all elements. Because we consider the photons only around the core regions, we ignore the absorption by torus.

The optical-UV accretion-disk spectral energy distributions (SEDs) are taken from Elvis et al. (1994). Here, the primary 2 keV X-ray disk luminosity is connected to the accretion-disk luminosity at 2500 Å

\[
\log L_{2 \text{ keV}} = 0.760 \log L_{2500 \text{ Å}} + 3.508 \tag{15}
\]

based on the study of 545 X-ray selected type 1 AGNs from the XMM-COSMOS survey (Lusso et al. 2010). Between UV and X-ray, following Lusso et al. (2010), we linearly connect the UV luminosity at 500 Å to the luminosity at 1 keV. Figure 1 shows the broadband AGN SED arising from the core region for various X-ray luminosities. AGN core SEDs typically have a spectral peak at \( \sim 30 \text{ eV} \), corresponding to \( \sim 10^8 \text{ K} \) (Figure 1), which corresponds to the emission radius at around \( \sim 10 R_g \).

![Figure 1. Typical broadband spectral energy distribution arising from the core region of AGNs. From top to bottom, each curve corresponds to 2–10 keV luminosity of \( 10^{46}, 10^{47}, 10^{48} \text{ erg s}^{-1} \), respectively.](image)

3.2. Physical Properties of Coronae

X-ray spectral studies allow us to determine some of the coronal parameters, such as the coronal electron temperature \( kT_e \) and the Thomson scattering optical depth \( \tau_T \) (e.g., Brenneman et al. 2014). We define \( k \) as the Boltzmann constant, and \( T_e \) is the electron temperature in Kelvin. The spectral cutoff at \( \sim 300 \text{ keV} \) of AGN core spectra corresponds to the electron temperature of \( kT_e \sim 100 \text{ keV} \). The process of Comptonization by thermal plasma is described by the Kompaneets equation (Kompaneets 1957). In this work, the photon index of the primary emission is assumed to be 1.9. This corresponds to \( \tau_T \sim 1.1 \), based on the solution to the Kompaneets equation (Zdziarski et al. 1996)

\[
\Gamma = \frac{9}{4} \theta_e^2 \left[ \tau_T (\tau_T + 1/3) \right] - \frac{1}{2}, \tag{16}
\]

where the dimensionless electron temperature \( \theta_e \equiv kT_e/m_e c^2 \). Therefore, in this work, we adopt \( kT_e = 100 \text{ keV} \) and \( \tau_T = 1.1 \). These values are consistent with the results from detailed X-ray spectral analysis (e.g., Fabian et al. 2015).

Recently, utilizing X-ray and radio data, Inoue & Doi (2018) found that the coronal magnetic field strength \( B \) is approximately 10 Gauss on scales of \( \sim 40 R_g \) from the SMBHs for two nearby Seyferts whose BH masses are \( \sim 10^8 M_\odot \). This coronal size is consistent with optical–X-ray spectral fitting studies (Jin et al. 2012) and microlensing observation (Morgan et al. 2012). Thus, in this paper, we set the coronal size as \( 40 R_g \) for all SMBHs and \( B = 10 \text{ G} \) for \( 10^8 M_\odot \) SMBHs.

Inoue & Doi (2018) also suggested that the coronae are likely to be advection-heated hot accretion flows (Kato et al. 2008; Yuan & Narayan 2014) rather than magnetically heated corona (Haardt & Maraschi 1991; Liu et al. 2002), because the measured magnetic field strength is too weak to keep the coronae hot and is rather consistent with the value based on the self-similar solutions of hot accretion flows (Kato et al. 2008; Yuan & Narayan 2014). Thus, we assume that coronal

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8 Contrary to this observational result, recent numerical simulations of the hot accretion flows (e.g., Kimura et al. 2019) show the magnetic field to be enhanced more by the magnetorotational instability (MRI; Balbus & Hawley 1991, 1998).
magnetic field strength scales as

\[ B \propto M_{\text{BH}}^{-1/2}, \]  

(17)

following the self-similar solution for the hot accretion flow (Yuan & Narayan 2014) where we ignore dependence on accretion rate and other parameters, for simplicity.

Mayers et al. (2018) have recently investigated a relation between the intrinsic 2–10 keV X-ray luminosity and the mass of central SMBHs using AGNs from the XMM-Newton Cluster Survey. The empirical relation found in Mayers et al. (2018) is given as

\[ M_{\text{BH}} = 2 \times 10^7 M_s \left[ \frac{L_{2-10 \text{ keV}}}{1.155 \times 10^{43} \text{ erg s}^{-1}} \right]^{0.746}. \]  

(18)

Using this relation, we convert X-ray luminosities to masses of central SMBHs.

### 3.3. Internal Gamma-Ray Attenuation in Coronae

Accelerated electrons and protons in coronae would emit gamma-rays (see Section 3.1). However, high-energy gamma-ray photons are attenuated by photon–photon pair production interactions \((\gamma \gamma \rightarrow e^+e^-)\) with low-energy photons. For isotropic target photons, the pair production cross section achieves its maximum of \(\approx 0.2 \sigma_T\) when a gamma-ray of energy \(E_{\gamma}\) interacts with a low-energy photon with energy (e.g., Aharonian 2004)

\[ \epsilon_{\text{peak}} \approx 3.5 m_e^2 c^4 \frac{E_{\gamma}}{E_{\gamma}} \approx 1 \left( \frac{1 \text{ TeV}}{E_{\gamma}} \right) \text{eV}. \]  

(19)

In terms of wavelength, \(\lambda_{\text{peak}} \approx 1.4(E_\gamma[\text{TeV}]) \mu\text{m}\).

Abundant photons are emitted from the AGN core region (Figure 1). From the SED of AGN core regions, as given in Section 3.1, we can compute the optical depth for high-energy gamma-rays to \(\gamma\gamma\) pair production interactions. The cross section for this process is (Breit & Wheeler 1934; Heitler 1954)

\[ \sigma_{\gamma\gamma}(E_{\gamma}, \epsilon, \theta) = \frac{3 \sigma_T}{16} (1 - \beta^2) \times \left[ 2 \beta (\beta^2 - 2) + (3 - \beta^4) \ln \left( \frac{1 + \beta}{1 - \beta} \right) \right], \]  

(20)

where \(\beta\) is

\[ \beta \equiv \sqrt{1 - \frac{2 m_e^2 c^4}{\epsilon E_\gamma (1 - \mu)}}; \quad \mu \equiv \cos \theta, \]  

(21)

where \(\theta\) is the angle between the colliding photons’ momenta.

For a photon with an energy of \(E_{\gamma}\), the \(\gamma\gamma\) optical depth is

\[ \tau_{\gamma\gamma}(E_{\gamma}) = \int_1^\infty d\mu \int_\epsilon_\text{in}^\infty \frac{1 - \mu}{2} \frac{U_{\text{ph}}(\epsilon)}{\epsilon^2} \sigma_{\gamma\gamma}(E_{\gamma}, \epsilon, \theta) \tau \]  

(22)

where \(\epsilon_{\text{in}}\) is the pair production threshold energy,

\[ \epsilon_{\text{in}} = \frac{2 m_e^2 c^4}{E_{\gamma} (1 - \mu)}. \]  

(23)

Integration over the interaction angle in Equation (22) can be performed analytically resulting in the angle averaged \(\gamma\gamma\) cross section (Aharonian 2004):

\[ \sigma_{\gamma\gamma} = \frac{3 \sigma_T}{2 \epsilon^2} \left[ \left( s + \frac{1}{2} \ln s - \frac{1}{6} + \frac{1}{2s} \right) \ln(\sqrt{s} + \sqrt{s - 1}) - \left( s + \frac{4}{9} \frac{1}{9s} \right) \sqrt{1 - s} \right], \]  

(24)

where \(s = E_{\gamma} c / m_e^2 c^4\).

Figure 2 shows the internal gamma-ray optical depth in the core region for various X-ray luminosities. The core region is expected to be optically thick against gamma-ray photons above 10–100 MeV, depending on disk luminosities. Such high optical thicknesses against pair production in AGN coronae are well-known (e.g., Bonometto & Rees 1971; Done & Fabian 1989; Fabian et al. 2015), based on the compactness parameter argument (Guilbert et al. 1983).

### 4. Timescales

Given the observed properties of AGN core regions, we can estimate the various timescales of high-energy particles in the coronae. Figure 3 shows the cooling rates of electrons in the coronae for different energy-loss processes, along with the acceleration rate and the freefall timescale following Section 2 and parameters presented in Section 3. We set \(\eta_{fi} = 30\) in the figure, which reproduces the IceCube neutrino background fluxes as discussed later in Section 7. The panels correspond to respective 2–10 keV X-ray luminosities of \(10^{42}, 10^{44},\) and \(10^{46}\) erg s\(^{-1}\).

Due to the intense broadband radiation field, the cooling is dominated by the Compton cooling. However, in higher energy regions, the main cooling channel is replaced by synchrotron cooling because of the KN effect. The more luminous AGNs tend to have a more efficient IC cooling effect, as the target photon density increases. When we assume \(\eta_{fi} = 30\), electron acceleration up to \(\gamma_e \sim 10^5\) (~50 GeV) is feasible in AGN coronae at various luminosities. Therefore, synchrotron radiation through coronal magnetic fields and gamma-ray emission by Comptonization of disk photons are naturally expected in AGN coronae.

ALMA spectra of two nearby Seyferts, whose X-ray luminosities are about \(10^{44}\) erg s\(^{-1}\), extends their radio
synchrotron power-law spectra at least up to 230 GHz, which corresponds to $\gamma_e \sim 80$ given the magnetic field strength of 10 G (Inoue & Doi 2018). As shown in the top right panel (the case of $\log L_X = 44$) in Figure 3, relativistic electrons with $\gamma_e \sim 80$ seen by ALMA can be easily accelerated in AGN coronae. Notably, such electrons can be accelerated even by a low-efficiency acceleration process, e.g., with $\eta_g \sim 10^6$. For this energy, Compton cooling is the dominant energy-loss process. As the cooling timescale for $\gamma_e \sim 80$ is about 100 s, flux variability in the radio synchrotron emission is expected; some Seyferts are already known to show a flux variation at least on day scales (Baldi et al. 2015). Further dense light curve observations may see shorter timescale variabilities.

Similar to Figure 3 for electrons, Figure 4 shows the timescales for high-energy protons for various luminosities. As in Figure 3, we set $\eta_g = 30$. Because synchrotron and Compton cooling are not effective for protons in our case, we do not show these timescales in the figure.

It is evident that protons can be accelerated up to $\gamma_p \sim 10^6$ ($\sim$1 PeV) in AGN coronae for various luminosities. The maximum attainable energy is controlled by different processes for different luminosity AGNs, due to SED and size dependence. For low-luminosity Seyferts ($L_X < 10^{44}$ erg s$^{-1}$), acceleration is limited by the dynamical timescale rather than radiative cooling, while it becomes limited by the Bethe–Heitler cooling for higher-luminosity objects. As the luminosity increases, $p\gamma$ and Bethe–Heitler cooling effects become more prominent. At higher luminosities, the Bethe–Heitler processes dominate the energy loss process of high-energy particles. Therefore, in cases of high-luminosity objects, the resulting hadronic gamma-ray and neutrino spectra in the TeV band will show spectral suppression due to the Bethe–Heitler processes; e.g., see Murase (2008), for cases of gamma-ray bursts.

5. Particle Spectrum

The steady-state particle distributions $n = dN/d\gamma$ can be derived from the solution of the transport equation (Ginzburg & Syrovatskii 1964)

$$\frac{\partial}{\partial \gamma} (\dot{\gamma}_{\text{cool}} n) + \frac{n}{t_{\text{fall}}} = Q(\gamma),$$

where $\dot{\gamma}_{\text{cool}}$ is the total cooling rate, and $Q(\gamma)$ is the injection function, which describes phenomenologically some acceleration process, e.g., DSA. The injection function for nonthermal protons and electrons is set as $Q(\gamma) = Q_0 \gamma^{-\gamma_{\text{inj}}} \exp(-\gamma/\gamma_{\text{max}})$. Here, $\gamma_{\text{max}}$ is the maximum Lorentz factor determined by balancing the acceleration and cooling timescales (Figures 3 and 4). The corresponding solution is

$$n = \left[\frac{1}{\dot{\gamma}_{\text{cool}}} \int_{\gamma_{\text{inj}}}^{\gamma_{\text{max}}} Q(\gamma') e^{-T(\gamma,\gamma')} d\gamma'\right]^{-1},$$
\[
T(\gamma_1, \gamma_2) = \frac{1}{t_{\text{fall}}} \int_{\gamma_1}^{\gamma_2} \frac{d\gamma}{\xi_{\text{cool}}}.
\]  

(27)

By solving Equation (26), we obtain a steady-state spectrum of the nonthermal particles.

Figure 5 shows the steady-state nonthermal electron spectrum obtained for the injection spectral index of \( p_{\text{inj}} = 2.0 \), along with the observationally determined electron spectral distribution for IC 4329A (Inoue & Doi 2018). ALMA observed nonthermal synchrotron radiation between 90.5 and 231 GHz, which corresponds to the electron Lorentz factors between 50 and 80, respectively. The corresponding region is shown as the shaded area in Figure 5.

For the calculation of the steady-state spectrum, we set \( M_{\text{BH}} = 10^8 M_\odot \), \( r_c = 40 \), \( B = 10 \text{ G} \), \( kT_e = 100 \text{ keV} \), \( \tau_T = 1.1 \), and \( \eta_g = 30 \). The synthetic electron distribution obtained for \( p_{\text{inj}} = 2.0 \) nicely reproduces the observationally determined electron spectrum in the energy range constrained by the observations. This injection index is naturally expected in a simple DSA scenario for a strong shock.

The resulting particle spectrum at \( \gamma_e > 10^4 \) becomes softer than the observationally determined index at \( 50 \lesssim \gamma_e \lesssim 80 \). This is because of the influence of the cutoff imposed by the particle cooling. Therefore, if we consider the high-energy synchrotron or IC spectral shapes, the cooling effects should be accurately taken into account. Even though the electron spectrum extends down to lower energies, it is hard to see the corresponding synchrotron emission due to synchrotron self-absorption effect (Inoue & Doi 2014).

The calculated electron spectrum is renormalized to agree with the observationally determined spectrum, which is achieved if the nonthermal electrons contains \( f_{\text{nth}} = 0.03 \) of the energy in thermal leptons. We note that, in order to define the energy content in the nonthermal particles, we formally integrate above \( \gamma_e = 1 \) in this study. We keep this fraction for nonthermal electron energy fixed in calculations below for all Seyferts.

The energy fraction of nonthermal electrons was fixed to \( \xi_{\text{nth}} = 0.04 \) in Inoue & Doi (2018). Here, \( \xi_{\text{nth}} \) is defined beyond the break electron Lorentz factor, while \( f_{\text{nth}} \) is above \( \gamma_e = 1 \).
That amount of nonthermal electrons overproduces the MeV background flux (see Section 7). To be consistent with the observed cosmic MeV gamma-ray background flux, we set $\xi_{\text{nth}} = 0.015$ in this work, which corresponds to $f_{\text{nth}} = 0.03$. The obtained best-fit parameters with this fraction for the radio spectrum of IC 4329A are $p = 2.9 \pm 0.9$, $B = 11.4 \pm 5.6$ G, and $r_c = 42.7 \pm 7.8$, which are very similar to those obtained for the case of $\xi_{\text{nth}} = 0.04$. We adopt these parameters for the observationally determined electron distribution in the Figure 5. Fitting results for the other parameters were also the same as those with $\xi_{\text{nth}} = 0.04$.

Here, the total shock power $P_{\text{sh}}$ can be estimated as

$$P_{\text{sh}} = 4\pi R_c^2 n_p m_p v_{\text{sh}}^3 / 2$$

$$\approx 2.2 \times 10^{\gamma} \left( \frac{\tau_r}{1.1} \right) \left( \frac{r_c}{40} \right)^{-1/2} \left( \frac{M_{\text{BH}}}{10^{8} M_{\odot}} \right) \text{erg s}^{-1}. \quad (28)$$

For objects with $L_X = 10^{44}$ erg s$^{-1}$, $f_{\text{nth}} = 0.03$ corresponds to $\sim 5\%$ of the shock power injected into the acceleration of electrons. This high value implies that, if DSA is responsible for particle acceleration in AGN coronae, then processes regulating injection of electrons into DSA are very efficient. For example, in the case of DSA in supernovae remnants, nonthermal electrons obtain only $\sim 1\%$ of energy transferred to nonthermal protons (Ackermann et al. 2013). Detailed consideration of the reasons for this unusually high efficiency of electron acceleration is beyond the scope of this paper, but we note that a significant presence of positrons may affect the ratio (e.g., Park et al. 2015). Given these uncertainties, we stipulate that the same energy injection rate is achieved for protons as for electrons. This power appears to be sufficient to explain the observed IceCube neutrino fluxes.

For the other object, NGC 985, the observed electron spectral index is $2.11 \pm 0.28$ (Inoue & Doi 2018), which is hard, considering the radiative cooling effect. Cascade components would have such a hard spectrum below the threshold energy (e.g., Aharonian & Plyasheshnikov 2003). In addition, due to the quality of data at low frequencies, we could not precisely determine the other components, such as free–free emission and synchrotron emission from star formation activity and synchrotron emission from the jet. Those uncertainties may have resulted in a less reliable measurement of the corona emission spectrum slope. Further observations are required to precisely determine the radio spectral properties in NGC 985.

6. Gamma-Rays and Neutrinos from AGN Coronae

Accelerated electrons and protons in AGN coronae generate gamma-ray and neutrino emission through IC scattering, $pp$ interaction, and $p\gamma$ interaction. Adopting a steady-state particle spectrum, we calculate the resulting gamma-ray and neutrino spectra from AGN coronae. We follow the method of Blumenthal & Gould (1970) to calculate the gamma-ray emission due to the IC scattering by nonthermal electrons. We calculate the gamma-ray and neutrino emission induced by hadronic interactions following Kelnner et al. (2006) for $pp$ interactions and Kelnner & Aharonian (2008) for $p\gamma$ interactions. For simplicity, we do not take into account IC-scattered emission from secondary electrons and positrons. For the thermal Comptonization spectra, we adopt the AGN SED shown in Figure 1, which takes into account reflection components but does not account for attenuation by torus.

The torus attenuation is mainly relevant for $\lesssim$30 keV, which is below the range of our interest.

Figure 6 shows the resulting gamma-ray and neutrino spectra for two cases. The neutrino flux is shown in the form of per flavor. The left panel of the figure shows the case with a 2–10 keV luminosity of $10^{43}$ erg s$^{-1}$ at a distance of 14 Mpc, while the right panel shows the case with a luminosity of $10^{44}$ erg s$^{-1}$ at a distance of 69 Mpc. The former and the latter roughly correspond to NGC 4151 and IC 4329A, respectively. NGC 4151 is the brightest Seyfert in the X-ray sky (Oh et al. 2018). For comparison, the overall fluxes of both panels are renormalized to match with the Swift/BAT flux of NGC 4151 and IC 4329A, respectively, at 14–195 keV (Oh et al. 2018). We note that we do not calculate the detailed X-ray spectra of each object, as doing so is beyond the scope of this paper.

We set the injection spectral index of $p_{\text{inj}} = 2.0$ and the gyrofactor of $\eta_{\text{g}} = 30$ for both electrons and protons (see Section 5). We also set the same injection power into protons and electrons as described in Section 5. The target photon density for IC scatterings and $p\gamma$ interactions is defined as $U_{\text{ad}}(\gamma)$ (see Section 3.1). Because we assume a uniform spherical source, gamma-ray photons are attenuated by the internal photon field by a factor of $3U_{\text{ad}}(\gamma)/\tau_{\text{int}}$, where $U_{\text{ad}}(\gamma) = 1/2 + \exp(-\gamma)/\gamma - [1 - \exp(-\gamma)]/\gamma$ (see Section 7.8 in Dermer & Menon (2009)), where $\tau_{\text{int}}$ is the internal gamma-ray optical depth (see Section 3.3). Gamma-rays are also attenuated by the extragalactic background light (EBL) during propagation in the intergalactic space. We adopt Inoue et al. (2013a) for the EBL attenuation.

For comparison, we also show the expected sensitivity curve of planned MeV missions: COSI-X (300 days), COSI-X (3 yr; De Angelis et al. 2017), e-ASTROGAM (35 days; Aramaki et al. 2019), and GRAMS (3 yr; Aramaki et al. 2019). The 10 yr sensitivity of Fermi/LAT is also shown. We also plot the sensitivity of neutrino detectors: IceCube and IceCube-Gen2 (van Santen & IceCube-Gen2 Collaboration 2017). We assume a decl. $\delta$ of 30$^\circ$ for the left panel, and $-30^\circ$ for the right panel.

Because the spectral index of electrons is $\sim$3 after radiative cooling, the resulting nonthermal gamma-ray spectrum is flat in $\nu F_{\nu}$ in the MeV band that appears after the thermal cutoff. Given the cooling–limited maximum energy $\gamma_e \sim 10^7$, the intrinsic IC spectrum can extend up to $\sim$100 GeV. However, due to the strong internal gamma-ray attenuation effect, the spectra will have a cutoff around 100 MeV in both cases. In the sub-MeV band, the spectra show superthermal tails due to the combination of thermal and nonthermal components and a spectral hardening at $\sim$1 MeV. These superthermal and flat spectral tails should be tested by future MeV gamma-ray missions. Balloon flights such as GRAMS (Aramaki et al. 2019) and SMILE (Takada et al. 2011; Komura et al. 2017) may be able to catch this superthermal tail, and satellite-class MeV
missions such as e-ASTROGAM (De Angelis et al. 2017), AMEGO,15 and GRAMS (Aramaki et al. 2019) will be able to see the nonthermal power-law tail as well. For the case of NGC 4151, Fermi/LAT may be able to see the signature with its 10 yr survey. However, the expected flux is almost at the sensitivity limit. Thus, it may need further exposures for Fermi/LAT to see the coronal emission.

The pp and pγ production efficiencies are given by the ratio between the dynamical timescale (Equation (2)) and the interaction timescales (Equations (8) and (10)). The pp production efficiency is analytically given as

$$f_{pp} = \frac{t_{\text{dynamical}}}{t_{pp}} \simeq 0.16 \left( \frac{r}{1.1} \right) \left( \frac{F}{40} \right)^{-1/2}.$$  

Gamma-rays and neutrinos induced by hadronic interactions carry one-third and one-sixth of those interacted hadron powers. Therefore, the hadronic gamma-ray and neutrino luminosities are expected to make up ~5% and ~3% of the intrinsic proton luminosity, respectively. Because we assume the same energy injection to electrons and protons, and the coronal Thomson scattering optical depth is 1.1, before the attenuation, we have hadronic gamma-ray and neutrino fluxes of ~5% and ~3% of the IC gamma-ray fluxes.

The pp- and pγ-induced gamma-rays are also mostly attenuated by the internal photon fields. Thus, we do not expect any GeV gamma-ray emission from Seyferts. Moreover, the intrinsic gamma-ray energy fluxes due to hadronic interactions are about a factor of 10 less than those due to primary electrons, because of radiative efficiency differences between protons and electrons. This implies that gamma-rays produced by secondary pairs should not significantly alter the resulting spectra. Therefore, we can safely ignore the cascade contribution.

Unlike the case of gamma-rays, neutrinos induced by hadronic interactions can escape from the system without any attenuation. Because we adopt the same p\text{max} = 2 for protons as for electrons, we expect a flat νFν spectrum for neutrinos, to which pp makes dominant contribution. At higher energies, especially in the case of IC 4329A, pp and pγ spectra are suppressed due to the Bethe–Heitler cooling process. The exact position of the cutoff energy depends on the assumed ηe. Here, as described later, we set ηe = 30 in order to be consistent with the IceCube background flux measurements. This gyrotropic results in a neutrino spectral cutoff around 100 TeV. Although it is difficult to see neutrino signals from individual Seyferts with the current generation of IceCube, it would be possible to see bright Seyferts in the northern hemisphere in the era of IceCube-Gen2 (see also Murase & Waxman (2016), for more general arguments). Therefore, even though Seyferts are faint in the GeV gamma-ray band, future MeV gamma-ray and TeV neutrino observations will allow our scenario to be tested.

7. Cosmic Gamma-Ray and Neutrino Background Fluxes From High-energy Particles in AGN Corona

In this section, we calculate the cosmic gamma-ray and neutrino background spectra from AGN coronae. For the cosmological evolution of AGNs, we follow Ueda et al. (2014), in which the evolutionary functions are defined at an intrinsic X-ray luminosity of 2–10 keV. We briefly review their formalism here.

Based on the luminosity-dependent density evolution model, the AGN X-ray luminosity function at a given luminosity L_X and a given redshift z is defined as

$$\frac{d\Phi_X(L_X, z)}{d \log L_X} = \frac{d\Phi_X(L_X, 0)}{d \log L_X} e^{-c(z, L_X)}.$$  

![Figure 6](image-url)
where \( d \Phi_X(L_X, 0)/d \log L_X \) is the luminosity function in the local universe, defined as

\[
d\Phi_X(L_X, z = 0) = A[(L_X/L_\ast)^\gamma + (L_X/L_\ast)^\gamma_z]^{-1},
\]

where \( A \) is the normalization and \( L_\ast \) is the break luminosity. Here, \( e(z, L_X) \) is the evolution factor represented as

\[
e(z, L_X) = \begin{cases} (1+z)^{p_1} & [z \leq z_c(L_X)], \\ (1+z_c)^{p_1}\left(1+z\right)^{p_2} & [z_c(L_X) < z \leq z_{c2}], \\ (1+z_c)^{p_1}\left(1+z_c\right)^{p_2}\left(1+z\right)^{p_3} & [z > z_{c2}]. \end{cases}
\]

The luminosity dependence for the \( p_1 \) parameter is considered as

\[
p_1(L_X) = p_1^* + \beta_1(\log L_X - \log L_p),
\]

where we set \( \log L_p = 44 \). Both cutoff redshifts are given by power-law functions of \( L_X \) as

\[
z_c(L_X) = \begin{cases} z_{c1}^*(L_X/L_{a1})^{p_1} & [L_X \leq L_{a1}], \\ z_{c1}^* & [L_X > L_{a1}], \end{cases}
\]

and

\[
z_{c2}(L_X) = \begin{cases} z_{c2}^*(L_X/L_{a2})^{p_2} & [L_X \leq L_{a2}], \\ z_{c2}^* & [L_X > L_{a2}], \end{cases}
\]

The parameters are summarized in Table 4 of Ueda et al. (2014). There is also a substantial fraction of Compton-thick AGNs in the universe (e.g., Ueda et al. 2003; Ricci et al. 2015). In order to take this population into account, we multiply the normalization factor by a factor of 1.5; for details, see Ueda et al. (2014).

The cosmic gamma-ray background fluxes are calculated as

\[
E^2 \frac{dN}{dE} = \frac{c}{4\pi} \int_{0.002}^{5} dz \int_{a_1}^{47} d\log L_X \left| \frac{dt}{dz} \right| \frac{d\Phi_X(L_X, z)}{d \log L_X} \times \frac{L_\gamma(E', L_X)}{1+z} \times \frac{3\tau_{\text{int}}(E', \log L_X)}{\tau_{\text{int}}(E', \log L_X)} \times \exp(-\tau_{\text{EBL}}[E, z]),
\]

where \( E' = (1+z)E \) and \( L_\gamma(E, L_X) \) represent the gamma-ray luminosity at energy \( E \) for a given X-ray luminosity of \( L_X \). The redshift and luminosity ranges are selected to be the same as in Ueda et al. (2014). Here, \( \tau_{\text{int}} \) and \( \tau_{\text{EBL}} \) represent the gamma-ray optical depth due to the internal photon field and the EBL, respectively. We do not consider the cascade gamma-ray photons (e.g., Inoue & Ioka 2012), because the gamma-ray energy fluxes due to hadronic interactions are already subdominant, compared to those attributed to primary electrons.

The neutrino background fluxes can be also calculated in the same manner, ignoring the gamma-ray attenuation terms and replacing \( L_\gamma(E, L_X) \) with \( L_\nu(E, L_X) \), where \( L_\nu(E, L_X) \) is the neutrino intensity at an energy of \( E \) for a given X-ray luminosity of \( L_X \).
Contribution of thermal and nonthermal electrons, respectively. Contribution of reflection is included in the thermal contribution. Cosmic X-ray and MeV gamma-ray background spectrum data from HEAO-1 A2 (Gruber et al. 1999), INTEGRAL (Churazov et al. 2007), HEAO-1 A4 (Kinzer et al. 1997), Swift-BAT (Ajello et al. 2008), SMM (Watanabe et al. 1997), Nagoya-Balloon (Fukada et al. 1975), and COMPTEL (Weidenspointner et al. 2000) are also shown in the figure.

Figure 7 shows the cosmic X-ray/gamma-ray and neutrino background spectra from AGN coronae, assuming the case of $p_{\text{inj}} = 2.0$ and $\eta_g = 30$. We also plot the observed background spectrum data from HEAO-1 A2 (Gruber et al. 1999), INTEGRAL (Churazov et al. 2007), HEAO-1 A4 (Kinzer et al. 1997), Swift-BAT (Ajello et al. 2008), SMM (Watanabe et al. 1997), Nagoya-Balloon (Fukada et al. 1975), COMPTEL (Weidenspointner et al. 2000), Fermi-LAT (Ackermann et al. 2015), and IceCube (Aartsen et al. 2015).

Figure 8 shows the cosmic X-ray background spectrum only from Figure 7. By setting $f_{\text{inh}} = 0.03$, the gamma-ray fluxes from AGNs coronae due to IC scattering by thermal and nonthermal electrons can nicely explain the observed cosmic MeV gamma-ray background radiation as an extension of the cosmic X-ray background radiation, which is known to be explained by Seyferts (Ueda et al. 2014). Because the spectral index of nonthermal electrons in the coronae is $\sim 3$, the resulting MeV gamma-ray background spectrum becomes flat in $E^2dN/dE$ (see Figure 8). Here, the cosmic X-ray background spectrum from Seyferts has a spectral cutoff above $\sim 300$ keV because of temperature of thermal electrons $\sim 100$ keV (Ueda et al. 2014). By summing up these two thermal and nonthermal components, superthermal tail appears in the sub-MeV band as observed by, e.g., Fukada et al. (1975), Kinzer et al. (1997), and Watanabe et al. (1997). Because the dominant IC contributors switches from thermal electrons to nonthermal electrons at around 1 MeV, the MeV background spectrum may have a spectral hardening feature at $\sim 1$ MeV. The result does not significantly change as far as $\eta_g < 1000$. If $\eta_g > 1000$, we may require lower $f_{\text{inh}}$.

Due to the internal gamma-ray attenuation effect, these nonthermal gamma-rays cannot contribute to the emission above GeV. For the same reason, most hadronic gamma-ray photons are attenuated by internal photon fields, resulting in the generation of multiple secondary particles. Because calculation of those populations lies beyond the scope of this paper, we ignore them in our estimate. Moreover, as we describe above, the intrinsic hadronic fluxes are already an order of magnitude below the leptonic fluxes. Thus, pairs induced by hadronic cascades will not significantly change our results.

Here, IC emissions due to nonthermal electrons also contribute in the X-ray band. Their contribution is about $\sim 5\%$ at 30 keV of the observed cosmic X-ray background flux, which may reduce the required number of the Compton-thick population of AGNs.

The model curve at $\lesssim 10$ keV slightly overproduces the measured background spectrum. This is because we do not take into account X-ray attenuation by torus. However, the treatment of those soft X-ray photons does not affect our results at all.

For neutrinos, the combination of $pp$ and $p\gamma$ interactions can nicely reproduce the IceCube fluxes below 100–300 TeV by assuming $\eta_g = 30$ and about 5% of the shock power into proton acceleration and the same for electrons. The $pp$ interactions dominate the flux at $\lesssim 10$ TeV, while $p\gamma$ interactions prevail above this energy. Because of the target photon field SED, $p\gamma$ is subdominant in the GeV–TeV band. If we inject more powers into protons, it inevitably overproduces the IceCube background fluxes. As $\gtrsim$ GeV gamma-rays are internally attenuated, AGN coronae emission will not be seen in GeV gamma-rays, even though they can make the IceCube neutrino fluxes. Such hidden cosmic-ray accelerators have been suggested as a possible origin of the IceCube neutrinos; see Murase et al. (2016), for a general argument.

Figure 9 shows the cosmic neutrino background spectra from AGN cores with various gyro factors ranging from 1 (Bohm limit) to $10^3$. It is clear that, if $\eta_g \ll 30$, the resulting neutrino fluxes overproduce the measured fluxes. On the other hand, if $\eta_g \gg 30$, AGN coronae cannot significantly contribute to the observed neutrino background fluxes. Thus, in order to explain the IceCube neutrino background fluxes by AGN cores, $\eta_g \sim 30$ is required. However, we note that these estimates are based on the assumed energy injection fraction to protons. Recent particle-in-cell simulations of proton–electron plasma considering radiatively inefficient accretion flows (RIAFs) showed that protons will carry several factors more energies than electrons (Zhdankin et al. 2018). If this is the case, larger $\eta_g$ is favored.
8. Discussion

8.1. Comparison with Previous Works on High-energy Neutrinos

In the literature, it has been argued that high-energy particles in the core of AGNs generate intense neutrino emission (e.g., Eichler 1979; Begelman et al. 1990; Stecker et al. 1992; Alvarez-Muñiz & Mészáros 2004). These originally predicted fluxes have been ruled out by high-energy neutrino observations (IceCube Collaboration 2005). However, recent studies have revisited the estimated fluxes and found that AGN core models can account for the entirety of the measured fluxes (Stecker 2013; Kalashev et al. 2015). In this section, we would like to compare our results with those recent studies (Stecker 2013; Kalashev et al. 2015).

The model suggest by Stecker (2013) is very similar to the originally proposed one (Stecker et al. 1992), but the background flux is assumed to be lower by a factor of 20. The original model was motivated by models explaining AGN X-ray spectra via the electromagnetic cascade emission of secondary particles (Kazanas & Ellison 1986; Zdziarski 1986), which is not the case based on current X-ray and gamma-ray observational results. The shock radius and the magnetic field strength was assumed to be 10R_s and 10^3 G in Stecker et al. (1992).

The model in Kalashev et al. (2015) is an extension of Stecker et al. (1992), adding the radial emission profile in the standard accretion disk into the consideration of the p_T cooling processes. In our modeling, we do not take such an anisotropic radiation field into account. However, given the observationally determined corona size, the dominant photon targets are likely to be generated in the inner region of the corona. The particle spectra in Kalashev et al. (2015) are fixed to match with the IceCube data.

Neutrino fluxes or cosmic-ray spectra are fixed to match with the latest IceCube data in Stecker (2013) and Kalashev et al. (2015). In this work, we take a more physical approach. Corona plasma density, corona size, and magnetic field strength are determined from observations (Inoue & Doi 2018) in our work. For example, we set R_c = 40 R_g and B = 10 G based on ALMA observations (Inoue & Doi 2018). With those parameters, we can follow the acceleration processes in coronae in the framework of DSA. We found that the AGN coronae can explain the IceCube neutrino background flux in the TeV band, if the gyrofactor is η_q = 30 and about 5% of the shock energy goes into proton acceleration. We also predict that next-generation MeV gamma-ray and neutrino experiments will be able to test our model by observing nearby bright Seyferts such as NGC 4151 and IC 4329A.

8.2. Plasma Condition in Corona

Considering the plasma density in the accreting corona, high-energy particles may have sufficient time to redistribute their kinetic energy through thermalization by elastic Coulomb (EC) collisions before the gas reaches the event horizon (Takahara & Kusunose 1985; Mahadevan & Quataert 1997). In this section, we discuss thermalization timescales of electrons and protons in the AGN coronae.

First, the electron thermalization timescale in the nonrelativistic regime is estimated to be (Spitzer 1962; Stepney 1983)

\[
\tau_{EC,ee} \approx \frac{4 \sqrt{\pi}}{n_e \sigma_T c \ln \Lambda} \frac{\theta_e^{3/2}}{2}
\]

\[
\approx 1.1 \times 10^3 \left(\frac{\tau_e}{1.1} \frac{r_e}{40} \frac{M_{BH}}{10^8 M_\odot}\right) \times \left(\frac{kT_e}{100 \text{ keV}}\right)^{3/2} \text{s},
\]

where \(\ln \Lambda \approx 20\) is the Coulomb logarithm. For relativistic electrons with Lorentz factors \(\gamma_e \gtrsim 1 + \theta_e\), the thermalization timescale due to interactions with the background plasma becomes (Dermer & Liang 1989)

\[
\tau_{EC,ee} (\gamma_e) = \frac{4}{3} \frac{K_2(\theta_e^{-1})}{n_e \sigma_T c (\ln \Lambda + 9/16 - \ln \sqrt{2})} \frac{\gamma_e^3}{\gamma_e - 1},
\]

where \(K_n\) is the modified Bessel function of order \(n\) and parameter \(u_{ee} = (\gamma_e^{-1} + \gamma_e^{-1} - 1 / 2) \theta_e\). This equation can be approximated as

\[
\tau_{EC,ee} (\gamma_e) \approx 2 \frac{\gamma_e}{3 n_e \sigma_T c (\ln \Lambda + 9/16 - \ln \sqrt{2})} \left[ \frac{K_2(\theta_e^{-1})}{K_2(\theta_e^{-1}) - 1} \right]^{-1}.
\]

This is a good analytic approximation at \(\theta_e \gtrsim 0.3\) and \(\gamma_e \gtrsim 2\) (Dermer & Liang 1989).

Second, the proton–proton relaxation timescale in the nonrelativistic regime is estimated to be (Spitzer 1962; Stepney 1983)

\[
\tau_{EC,pp} \approx \frac{4 \sqrt{\pi}}{n_p \sigma_p c \ln \Lambda} \frac{m_p^2}{m_e} \frac{\theta_p^{3/2}}{2}
\]

\[
\approx 4.7 \times 10^4 \left(\frac{\tau_e}{1.1} \frac{r_e}{40} \frac{M_{BH}}{10^8 M_\odot}\right) \times \left(\frac{kT_p}{100 \text{ keV}}\right)^{3/2} \text{s},
\]

where \(\theta_p = kT_p / m_p c^2\) is the dimensionless proton temperature. At high kinetic energies, nuclear interaction becomes important; see Gould (1982), for details. In the mildly relativistic case, the elastic proton–proton relaxation timescale approximately becomes (Gould 1982)

\[
\tau_{EC,pp} \approx \frac{4}{n_p \sigma_p c} \frac{\beta_p \gamma_p^2}{\gamma_p^2 - 1},
\]

where \(\sigma_t \approx 2.3 \times 10^{-26} \text{cm}^2\). This approximation is valid at 70 MeV \(\lesssim (\gamma - 1) m_e c^2 \lesssim 500 \text{ MeV}\). Above 500 MeV, inelastic processes start to dominate.

Finally, the proton–electron thermalization timescale due to EC collisions in the nonrelativistic regime is estimated to be
Figure 10. Electron and proton thermalization timescales in AGN coronae, together with radiative cooling and dynamical timescales. Thick solid curve shows the freefall timescale. Dashed, dotted, and dot-dashed curves correspond to synchrotron cooling, IC cooling, and ee EC thermalization timescale for electrons, respectively. Double-dotted-dashed, triple-dotted-dashed, and thin solid curves correspond to pp EC thermalization, pe EC thermalization, and pp inelastic interaction timescale for protons, respectively. We set log \( L_X = 44 \), \( \tau_T = 1.1 \), \( R_t = 40 \) \( R_s \), and \( kT_e = kT_p = 100 \) keV.

(Spitzer 1962; Stepney 1983)

\[
t_{\text{EC,pp}} \approx \frac{\sqrt{\pi/2}}{n_e \sigma_T e \ln\Lambda} \left( \frac{m_p}{m_e} \right) (\theta_p + \theta_p)^{3/2} \\
\gtrsim 3.6 \times 10^5 \left( \frac{\tau_T}{1.1} \right)^{-1} \left( \frac{r_e}{40} \right) \left( \frac{M_{\text{BH}}}{10^8M_\odot} \right) \\
\times \left( \frac{kT_p}{100 \text{ keV}} \right)^{3/2} \text{(s)},
\]

where we assume \( \theta_p = \theta_p \). The temperature of a hot accretion can roughly reach virial temperature \( T_p \approx GM_{\text{BH}} m_p/3kR \sim 3 \times 10^{32}(R/R_s)^{-1} \) K. At such a higher temperature, \( t_{\text{EC,pp}} \) becomes longer. In the case of relativistic protons, the energy-loss timescale through elastic Coulomb (EC) interactions is given as (Mannheim & Schlickeiser 1994; Dermer et al. 1996)

\[
t_{\text{EC,pp}} \approx 1.2 \times 10^6 \left( \frac{3(\beta_p^2 + \beta_p^3)(\gamma_p - 1)}{n_e \sigma_T c \beta_p^2 \ln\Lambda} \right),
\]

where \( \beta_p = \sqrt{1 - 1/\gamma_p^2} \). At \( \gamma_p \gg 1 \) and \( \theta_p \ll 1 \), the relativistic EC scattering relaxation time can be approximated as

\[
t_{\text{EC,pp}} \approx 2.9 \times 10^5 \left( \frac{\tau_T}{1.1} \right)^{-1} \left( \frac{r_e}{40} \right) \left( \frac{M_{\text{BH}}}{10^8M_\odot} \right) \left( \frac{\gamma_p}{100} \right) \text{(s)}.
\]

Figure 10 shows EC thermalization timescales for electrons and protons for the luminosity of \( L_X = 10^{44} \) erg s\(^{-1}\). Because EC thermalization is effective for low-energy particles, the horizontal axis is shown in \( \gamma/\beta \).

Around \( \gamma/\beta \approx 2 \), \( t_{\text{EC,ee}} \) shows a sharp feature, which is related to the temperature of the background plasma, \( kT_e \approx 100 \) keV. At this temperature, the electron distribution has a peak around \( \sim 3kT_e \) corresponding to \( \gamma/\beta \approx 1.2 \). Thus, around this energy, mean energy transfer is small. We note that, below this energy, electrons gain energies from the background plasma through elastic ee scatterings rather than losing their energies (Dermer & Liang 1989), however, this energy gain process is not considered in our work, because it is not relevant for our energy range of interest. As seen in Figure 10, the energy-loss process of electrons is dominated by the Compton cooling at \( \gamma/\beta \gtrsim 1 \).

Following Gould (1982), we calculate the elastic pp timescale in the mildly relativistic regime. Because it assumes that an incident proton has much higher kinetic energy than background plasma, we combine the nonrelativistic \( t_{\text{EC,pp}} \) (Equation (40)) and that from Gould (1982). As discussed above, inelastic processes start to dominate at the kinetic energies of \( \gtrsim 500 \) MeV \( (\gamma_p \beta_p \gtrsim 1.2) \). For comparison, we also show inelastic pp interaction timescale \( t_{\text{pp}} \).

As the proton–electron Coulomb timescale \( t_{\text{EC,pe}} \) is longer than \( t_{\text{inel}} \), protons and electrons may not be in the thermal equilibrium in AGN coronae. The proton temperature of a hot accretion can roughly reach virial temperature \( T_p \approx GM_{\text{BH}} m_p/3kR \sim 3 \times 10^{32}(R/R_s)^{-1} \) K, which is \( \gg T_e \). Further, the existence of pairs in coronae can reduce \( n_p \).

Moreover, the shock-heated proton temperature becomes \( kT_p \sim 3m_p \gamma_p^2 \beta_p^2 \sim 4(r_e/40)^{-1} \) MeV. Those shock-heated protons and electrons also gain and lose their energies through these processes and would contribute as a thermal population in the coronae. These electrons are heated and cooled through EC proton–electron thermalization and Comptonization, respectively (e.g., Katz et al. 2011; Murase et al. 2011). The heating rate can be written as

\[
\frac{dT_p}{dt} = \frac{dT_p}{dt_{\text{EC,pe}}} = \frac{n_e \sigma_T e \ln\Lambda \left( \frac{m_e}{m_p} \right) T_p \theta_p^{3/2}}{\sqrt{\pi/2}},
\]

assuming \( \theta_p \gg \theta_p \). The cooling rate through Comptonization is

\[
\frac{dT_e}{dt} \approx -\frac{4 \sigma_T U_{\text{ph,tot}} T_e}{3 m_e c}.
\]

By equating these two heating and cooling rates of thermal electrons, the shock-heating electron temperature is estimated to be

\[
kT_e \approx k \left( \frac{3 \ln\Lambda m_e n_e}{4 \sqrt{\pi/2} m_p U_{\text{ph,tot}}} \right)^{2/5} T_p^{2/5}
\]

\[
\approx 86 \left( \frac{T_e}{1.1} \right)^{2/5} \text{(keV)},
\]

where we assume \( L_{\text{ph,bot}} \propto M_{\text{BH}} \). This temperature is close to the measured coronal temperature. Therefore, such a shock heating mechanism may be able to explain the current observed coronal temperature. To understand the detailed nature of thermal coronae, further studies including thermodynamical processes are required.

8.3. Other Particle Acceleration Mechanisms

In this paper, we consider the DSA as fiducial acceleration mechanism. However, other acceleration mechanisms such as turbulent acceleration, magnetosphere acceleration, and magnetic reconnection can also operate in AGN coronae. We briefly discuss these processes here.

First, turbulent acceleration is considered for low-accretion-rate objects, such as low-luminosity AGNs (e.g., Kimura et al. 2015; Zhdankin et al. 2017, 2018; Wong et al. 2019). In this scenario, particles are accelerated stochastically by turbulence and magnetic reconnection in accretion disk or coronae. Recently, Zhdankin et al. (2018) investigated electron–ion
plasma energization via turbulent dissipation in RIAFs using particle-in-cell simulations for the ion temperature $T_i$ in the range of $m_e c^2 \lesssim k_B T_i \lesssim m_p c^2$. Turbulent electron–ion plasma driven by MRIs generates power-law spectra for both species, and the indices depends on the initial ion temperature. The fractions of the kinetic energy in the nonthermal ions and electrons are $\sim 60\%$ and $6\%$ for ions and electrons at $k_B T_i \sim m_e c^2$, respectively. The fraction in nonthermal electrons is close to the required value for the MeV background (see Section 7).

We briefly follow the stochastic acceleration in the AGN coronae case. According to the quasilinear theory, the diffusion coefficient in the momentum space is (e.g., Dermer et al. 1996)

$$D_p \simeq (m_p c)^2 (ck_{\text{min}}) \left( \frac{v_A}{c} \right)^2 \zeta (r_L k_{\text{min}})^{q-2} \gamma^q, \quad (48)$$

where $k_{\text{min}} \sim R_s^{-1}$ is the minimum wave number of turbulence spectrum (corresponding to the size of the corona), $v_A = B/\sqrt{4\pi n_p m_p}$ is the Alfvén speed, $n_\ell = m_e c^2/eB$ is the Larmor radius, and $\zeta = \delta B^2 / B^2$ is the ratio of strength of turbulence fields against the background. The acceleration timescale is then estimated to be

$$t_{\text{ac}} \simeq \frac{p^2}{D_p} \simeq \frac{1}{\zeta} \left( \frac{v_A}{c} \right)^{-2} R \left( \frac{r}{x} \right)^{2-q} \gamma^{2-q}. \quad (49)$$

Assuming the Kolomogorov spectrum for the turbulent $(q = 5/3)$ and $\zeta = 1$, the timescale becomes

$$t_{\text{ac}} \simeq 3.1 \times 10^6 \left( \frac{r}{10} \right) \left( \frac{\gamma p}{100} \right)^{1/3} \left( \frac{B}{10 \text{ G}} \right)^{-7/3} \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{-1/3} \left( \frac{\gamma^2}{125} \right) \left( \frac{t_T}{1.1} \right)^{-1/2} \left( \frac{r}{40} \right)^{5/2} \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{5/2} \left( \frac{\gamma^2}{125} \right). \quad (50)$$

Thus, stochastic acceleration appears to be inefficient as compared to the typical cooling rates. This is caused by the measured weak magnetic fields, which result in small Alfvén speed. If the magnetic fields are amplified by MRIs, more efficient acceleration can be realized (e.g., Zhidanik et al. 2018).16

Second, magnetosphere acceleration can also accelerate particles in the vicinity of SMBHs (e.g., Beskin et al. 1992; Levinson 2000; Neronov & Aharonian 2007; Levinson & Rieger 2011; Rieger 2011). At low accretion rates, the injection of charges into the BH magnetosphere is not sufficient for a full screening of the electric field induced by the rotation of the compact object. The regions with unscreened electric field, so-called gaps, are able to effectively accelerate charged particles.

In order to have gaps, the maximum allowed accretion rate is (Levinson & Rieger 2011; Aleksić et al. 2014; Aharonian et al. 2017)

$$\dot{m} < 3 \times 10^{-4} \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{-1/7}, \quad (51)$$

where $\dot{m}$ is the accretion rate in the Eddington units. Because we are considering the standard accretion disk regime $\dot{m} \gtrsim 0.01$, particle acceleration by gaps will not occur in our case.

Finally, magnetic reconnection would accelerate particles; see, e.g., Hoshino & Lyubarsky (2012) for reviews. Reconnection would naturally happen in coronae as they are magnetized, and radiative magnetic reconnection has been suggested as an origin of the X-ray emission seen in accreting black hole systems (Beloborodov 2017). However, even in the case of solar flares, particle acceleration mechanisms in magnetic reconnection is still uncertain (e.g., Liu et al. 2008; Nishizuka & Shibata 2013). Although quantitative discussion is not easy here, the available energy injection power can estimated as

$$P_B = \frac{B^2 R_c^2 v_A}{2} \approx 5.4 \times 10^{39} \left( \frac{r}{1.1} \right)^{-1/2} \left( \frac{r}{40} \right)^{5/2} \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{5/2} \times \left( \frac{B}{10 \text{ G}} \right)^3 \text{ (erg s}^{-1}) \). \quad (52)$$

This power is not sufficient to provide the nonthermal particle energies. For detailed estimation, we may need to consider the spatial distribution of the magnetic field. However, such information is not currently available.

### 8.4. Cosmic MeV Gamma-Ray Background Radiation

It is known that Seyferts generate the cosmic X-ray background radiation (Ueda et al. 2014). The cosmic gamma-ray background at 0.1–820 GeV is believed to be explained by three components: blazars (e.g., Inoue & Totani 2009; Ajello et al. 2015), radio galaxies (Inoue 2011), and star-forming galaxies (Ackermann et al. 2012a). Unlike cosmic X-ray and GeV background radiation, the origin of the cosmic MeV gamma-ray background radiation is still veiled in mystery.

Nonthermal IC emission from coronae in Seyferts has been suggested as a possible scenario (Inoue et al. 2008). The MeV tail extended from the X-ray background spectrum is generated by nonthermal electrons with a very soft spectral index (Inoue et al. 2008). However, nonthermal electrons are included in an ad hoc way. In our work, we consider the particle acceleration and cooling processes given the latest observations. The tail is due to the superposition of thermal Comptonization cut-off spectrum and $\gamma\gamma$ attenuated flat nonthermal IC component. We can distinguish these two scenarios by observing individual objects in radio and X-ray bands.

Blazars are also considered potential candidates for the origin of the MeV background (Ajello et al. 2009). In order to distinguish Seyferts and blazars, we need to resolve the MeV sky. However, it is difficult to do this, and will remain so even with future MeV instruments (Inoue et al. 2015). We suggest that anisotropy measurements may distinguish these two scenarios (Inoue et al. 2013b) because blazar background should feature stronger Poisson fluctuations. Future MeV gamma-ray anisotropy observations will be important to

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16 After we submitted our paper to the journal and arXiv, a similar study on AGN coronae by Murase et al. (2019) appeared on arXiv. Both studies are independent and the greatest point of difference is the assumed particle acceleration processes. In our paper, we consider DSA, while Murase et al. (2019) consider stochastic acceleration motivated by recent numerical simulations (Kimura et al. 2019). However, as we discussed in this section, stochastic acceleration may not work, given the ALMA results of weak coronal magnetic field.
understand the particle acceleration in coronae and the origin of the MeV gamma-ray background radiation.

8.5. Gamma-ray Observations toward Seyferts

Gamma-rays from Seyfert galaxies are not yet robustly detected (Lin et al. 1993; Teng et al. 2011; Ackermann et al. 2012b). Possible signatures of gamma-ray emission above 0.1 GeV have been reported for ESO 17. The Athena X-ray observatory website (https://www.the-athena-x-ray-observatory.eu) reports that the coronae of Seyfert galaxies are about 10^{43} erg s^{-1}. The required luminosity ratio between X-ray and gamma-ray L_{0.1–10 GeV}/L_{14–195 keV} for these sources is about 0.1 (Ackermann et al. 2012b). Our model estimates this ratio as ~0.01. Therefore, coronal gamma-ray emission is most likely unable to account for the observed gamma-ray fluxes from those Seyfert galaxies. We note that the details would strongly depend on the properties of each object.

Although gamma-rays from other Seyferts have not been detected yet, Fermi/LAT has set upper limits on their gamma-ray fluxes (Teng et al. 2011; Ackermann et al. 2012b). Based on an analysis of the first 2–3 yr data, L_{0.1–10 GeV}/L_{14–195 keV} < 0.1 at the 95% confidence level has been obtained in most cases, which is consistent with our model estimate. The most stringent observational constraint is derived for NGC 4151, in which L_{0.1–10 GeV}/L_{14–195 keV} < 0.0025, even though the limit can vary with an assumed spectral shape. Following our models, the current 10 yr survey data of Fermi/LAT may be able to see NGC 4151 (Figure 1), even though the expected flux is almost at the sensitivity limit.

8.6. Fraction of Nonthermal Electrons

We set the energy fraction of nonthermal electrons in AGN coronae as f_{nth} = 0.03 because it nicely reproduces the observed MeV gamma-ray background radiation. As discussed in Inoue & Doi (2018), f_{nth}, B, and R_c are closely tied; current radio and X-ray data do not allow us to solve these three parameters simultaneously without decoupling thermal and nonthermal components.

Observationally, f_{nth} is constrained as <0.3 in order not to violate X-ray data based on NuSTAR observations (Fabian et al. 2017). If f_{nth} is significantly lower, it becomes difficult to attribute the MeV gamma-ray background radiation solely to Seyferts. However, a much lower value of f_{nth} contradicts other observations, because it requires a larger R_c based on the radio spectral fitting. If we set f_{nth} = 10^{-5} and 10^{-4}, R_c becomes ~70 R_g and ~100 R_g, respectively. The size of the corona is also constrained to an order of ~10 R_g by optical–X-ray spectral fitting studies (Jin et al. 2012) and microlensing observations (Morgan et al. 2012). Therefore, f_{nth} cannot become much smaller than the adopted value.

8.7. Nuclear Spallation in AGNs

Given the ALMA results, particle accelerations occur in AGN coronae. As we demonstrated, high-energy protons are easily accelerated in coronae. These high-energy protons can be also traced by future high-resolution calorimeter spectroscopy in the X-ray band, e.g., XRISM (Tashiro et al. 2018) and Athena (Nandra et al. 2013). As narrow line features are seen in AGN X-ray disk spectra, there must be abundant metal elements in AGN cores. Accelerated protons interact with those nuclei and induce nuclear spallation. The nuclear spallation in AGN disks will result in enhancement of emission lines from Mn, Cr, V, and Ti (Gallo et al. 2019). Those signatures will provide clues for the test of our model.

9. Conclusion

Recently, Inoue & Doi (2018) reported that the coronae of Seyferts are composed of both thermal and nonthermal electrons, based on ALMA observations, which implies that particle acceleration occurs in AGN coronae. In order to investigate the production mechanism of those high-energy particles, we study the particle acceleration process in AGN coronae. We consider particle acceleration by the DSA process in the coronae as an example. By taking into account the observationally determined coronal properties, such as temperature, density, size, and magnetic field strength, we find that standard DSA processes can easily reproduce the observed nonthermal electron in the coronae with an injection electron spectral index of p_{inj} = 2. Even in cases with low acceleration efficiency (\eta_p \sim 10^3), such populations can be realized in coronae. Given the observed magnetic field strength of 10 G and accretion rates, we also find that other possible acceleration mechanisms such as turbulent acceleration, magnetosphere acceleration, and magnetic reconnection confront difficulty in reproducing the observed nonthermal electrons.

The accelerated nonthermal electron populations will generate a MeV gamma-ray power-law spectrum in the AGN SEDs up to ~0.1 GeV, which is limited by internal gamma-ray attenuation. In the sub-MeV band, the spectrum shows a superthermal tail due to the combination of thermal and nonthermal components and spectral flattening occurs at ~1 MeV. These superthermal and flat spectral tails should be tested by future MeV gamma-ray missions.

We also study the contribution of AGN coronae to the cosmic gamma-ray background radiation. By setting the energy fraction of nonthermal electrons f_{nth} ~ 3%, corresponding to ~5% of the shock energy in electron acceleration, AGN coronae can explain the MeV background in an extension of the X-ray background contribution of Seyferts. Due to a strong internal gamma-ray attenuation effect, the contribution of AGN coronae to the GeV background is negligible.

Accelerated particles would also result in neutrino production through hadronic processes. Intense neutrino emission has been expected to be produced in AGN coronae once hadrons are accelerated together (e.g., Begelman et al. 1990; Stecker et al. 1992; Alvarez-Muñiz & Mészáros 2004). Recent studies have proposed that these AGN core models could reproduce the high-energy neutrino flux measured by IceCube (Stecker 2005, 2013; Kalashev et al. 2015). However, normalization of neutrino fluxes from AGNs and acceleration properties of high-energy particles in those models are assumed to match with the observation.

We found that AGN coronae can explain the diffuse neutrino fluxes below 100–300 TeV under specific parameters of energy injection rates in protons and gyro factors. The allowed parameter regions are quite narrow. Protons and electrons should have the same energy injection rate and the gyro factor \eta_p should be ~30. IceCube Gen-2 will be able to test this scenario by searching the neutrino signal from nearby Seyfert galaxies such as NGC 4151 and IC 4329A.
In summary, Seyfert coronae are feasible sites for particle acceleration. If the energy injection rate is $5\%$ for both protons and electrons and the gyro factor is $\eta_g = 30$, they may be able to simultaneously explain the cosmic X-ray, MeV gamma-ray, and TeV neutrino background radiation. Future MeV gamma-ray and TeV neutrino observations will be able to test this scenario via observations of nearby bright Seyferts.

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