VARIABLE TeV EMISSION AS A MANIFESTATION OF JET FORMATION IN M87?

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ABSTRACT

It is proposed that the variable TeV emission observed in M87 may be produced in a starved magnetospheric region, above which the outflow associated with the VLBA jet is established. It is shown that annihilation of MeV photons emitted by the radiative inefficient flow in the vicinity of the black hole can lead to injection of seed charges on open magnetic field lines, with a density that depends sensitively on accretion rate, \( n_{\perp} \propto \dot{m}^3 \). For an accretion rate that corresponds to the inferred jet power, and to a fit of the observed spectral energy distribution by an ADAF model, the density of injected pairs is found to be smaller than the Goldreich–Julian (GJ) density by a factor of a few. It is also shown that inverse Compton scattering of ambient photons by electrons (positrons) accelerating in the gap can lead to a large multiplicity, \( \sim 10^3 \), while still allowing photons at energies of up to a few TeV to freely escape the system. The estimated gap width is not smaller than 0.01\( r_g \) if the density of seed charges is below the GJ value. The very high energy power radiated by the gap can easily account for the luminosity of the TeV source detected by H.E.S.S. The strong dependence of injected pair density on accretion rate should render the gap emission highly intermittent. We also discuss briefly the application of this mechanism to Sgr A*.

Key words: galaxies: individual (M87) – gamma rays: galaxies – radiation mechanisms: non-thermal

1. INTRODUCTION

Combined Very Long Baseline Array (VLBA) and TeV observations of M87 reveal a rapidly varying TeV emission that appears to be associated with the m.a.s VLBA jet (Acciari et al. 2009). The rapid flaring activity of the TeV source, with timescales \( t = 1 \text{ day} \) as low as 1–2 days, implies a source size of \( d \sim 4 r_g \text{ day} \) for a black hole mass \( M_{\text{BH}} = 4 \times 10^9 \) solar masses.

This, and the fact that the TeV emission appears to be correlated with the VLBA jet but not with emission from larger scales (and in particular \( HST-I \)), motivates the consideration that the observed TeV photons originate from the black hole magnetosphere (Neronov & Aharonian 2007; Rieger & Aharonian 2008). A plausible magnetospheric process discussed in the literature is curvature and/or inverse Compton (IC) emission by particles, either hadrons or leptons, accelerating in a vacuum gap of a starved magnetosphere. An alternative explanation for the observed fast TeV variability is emission from small regions located at larger radii, \( r \sim 100 \text{ } r_p \), as, e.g., in the misaligned minijet model of Giannios et al. (2010), or perhaps interaction of the jet with red giant stars closer to its base (Barkov et al. 2010).

The presence of the VLBA jet implies that a force-free (or ideal MHD) flow is established on scales <100\( r_p \) (Walker et al. 2008; Acciari et al. 2009), so that the magnetosphere is anticipated to be screened in the sense that the invariant \( E \cdot B \) nearly vanishes everywhere. However, as in pulsar theory, there must be a plasma source that replenishes charges which escape the system (both to infinity and across the horizon) along the open magnetic field lines in the polar region. The nature of this plasma source is poorly understood at present.

The injection of charges into the magnetosphere may be associated with the accretion process. Direct feeding seems unlikely, as charged particles would have to cross magnetic field lines on a timescale shorter than the accretion time in order to reach the polar outflow. Free neutrons that may be produced in a radiative inefficient accretion flow (RIAF) can cross field lines, however, they will decay over a distance \( \sim 0.03 r_g \) for a \( 10^8 \text{ } M_\odot \) black hole and, even if existent at sufficient quantity, will not reach the inner regions. On the other hand, MeV photons that are emitted by the hot gas near the horizon can annihilate in the polar region to produce charged leptons. Below, it is shown that the density of the charges thereby injected depends sensitively on the accretion rate and the conditions in the RIAF. Naive estimates suggest that in the case of M87 this process cannot provide complete screening at accretion rates that correspond to the inferred jet power, and to a fit of the observed spectral energy distribution (SED) by an ADAF model. Those estimates are, however, highly uncertain, as explained below.

Another plausible plasma source is cascade formation in starved magnetospheric regions. The size of the gap then depends on the conditions in the magnetosphere and the pair-production opacity. As shown elsewhere (Levinson 2000), vacuum breakdown by back-reaction is unlikely, as it requires magnetic field strength in excess of a few times \( 10^5 \text{G} \), higher than the equipartition value for Eddington accretion. Pair production via absorption of TeV photons by the ambient radiation field is more likely. However, the spectrum of the very high energy (VHE) photons observed by H.E.S.S. extends up to \( \sim 10 \text{ TeV} \) (e.g., Aharonian et al. 2006), and the assumption that these photons originate from the magnetosphere (or even the VLBA jet) implies that the pair-production opacity at these energies must not exceed unity. This raises the question whether pair cascades in the magnetosphere can at all account for the multiplicity required to establish a force-free flow.

In what follows it is shown that IC scattering of ambient photons by electrons (positrons) accelerating in the gap can lead to a large multiplicity, \( \sim 10^3 \), while still allowing photons...
H. E. S. S., and the fluctuations of the resulted force-free outflow, as indicated by outflow is established just above the pair formation front and appears as the pair cascades just above the gap, leading to a large multiplicity. A force-free energies below a few TeV can escape freely to infinity. Interactions of IC

Figure 1. Schematic representation of the magnetosphere structure: a vacuum gap of height \( h < r_s \) accelerates particles (electrons or positrons) to high Lorentz factors. The gap is exposed to soft radiation emitted from a source of size \( R_g \). Curvature emission and inverse Compton scattering of ambient radiation produce VHE photons with a spectrum extending up to \( 10^4 \) TeV. Photons having force-free flow, as observed. A schematic illustration of the gap of height \( \beta \).

of both, the magnetospheric TeV emission and the resultant cascade formation process would naturally lead to the variability luminosity observed by H. E. S. S., and the fluctuations of the resulted force-free outflow, as indicated by the morphological changes of the VLBA jet.

at energies of up to a few TeV to freely escape the system. The electromagnetic cascade is initiated by IC photons having much higher energies, up to \( \sim 10^4 \) TeV, for which the \( \gamma \gamma \) -optical depth is much larger. The seed charges are provided by annihilation of MeV photons from the RIAF. It is found that the gap width is not smaller than 0.01\( r_s \) if the density of seed charges is below the Goldreich–Julian (GJ) value. The luminosity of the VHE photons produced in the gap can account for the TeV luminosity observed by H. E. S. S.

of the gas, and one can safely assume \( n_e \) in the logarithmic term in Equations (3.6) and (3.8) of NY95. Since the RIAF is optically thin, pair production does not affect the leptonic content of the gas, and one can safely assume \( n_e = n_i \). By employing Equation (1) one arrives at

\[
q_{\text{ff}} \simeq 1.8 \times 10^2 \theta_\gamma \dot{m} M_9^{-2} (r/r_s)^{-3} \text{ erg s}^{-1} \text{ cm}^{-3}. \tag{4}
\]

The free–free luminosity emitted by the RIAF is \( L_{\gamma \gamma} = \int q_{\gamma \gamma} d^3 r \simeq 2\pi r^3 q_{\gamma \gamma} \ln(r/r_s) \), from which we readily obtain the number density of MeV photons in the magnetosphere:

\[
n_{\gamma} = \frac{q_{\gamma \gamma} 2\pi r^3 \ln(r/r_s)}{2\pi c r^2 \epsilon_\gamma} \simeq 0.2 q_{\text{ff}} r^3 \epsilon_\gamma \simeq 1.4 \times 10^{11} \dot{m} M_9^{-1}. \tag{5}
\]

where \( \epsilon_\gamma = 3\theta_\gamma (m_e c^2) \). The production rate of \( e^\pm \) pairs inside the magnetosphere due to \( \gamma \gamma \)-annihilation is approximately \( \sigma_{\gamma \gamma} n_e^2 c (4\pi/3) r_s^2 \). In steady state this rate is balanced by the escape rate, roughly \( 4\pi r^2 \dot{m}_{\pm} c \). Equating the two rates one has

\[
n_{\pm} = \sigma_{\gamma \gamma} n_e^2 r_s / 3 \simeq 3 \times 10^{11} \dot{m} M_9^{-1} \text{ cm}^{-3}. \tag{6}
\]

Henceforth, the accretion rate is measured in units of the Eddington rate, \( \dot{m} = \dot{M} / \dot{M}_{\text{Edd}} \), where the Eddington accretion rate is defined as \( \dot{M}_{\text{Edd}} = L_{\text{Edd}} / n_{\text{ff}} c^2 = 10^7 M_9 \) \( \text{g s}^{-1} \), with \( n_{\text{ff}} = 0.1 \) adopted. The ion temperature of the accreted gas is close to virial, reaching \( T_i \sim 10^{12} \) K at \( r = r_s \). The ion density is given by

\[
n_i(r) = \frac{M}{2\pi r^2 m_p v_f} = 5 \times 10^{11} \dot{m} M_9^{-1} (r/r_s)^{-3/2} \text{ cm}^{-3}, \tag{1}
\]

adopting \( v_f = 0.1 c (r/r_s)^{-1/2} \) for the RIAF radial velocity. The Thomson optical depth is

\[
\tau \simeq 0.2 \dot{m} (r/r_s)^{-1/2}. \tag{2}
\]

Typically \( \tau \ll 1 \) for accretion rates \( \dot{m} \ll \dot{m}_{\text{crit}} \), where \( \dot{m}_{\text{crit}} \) is the critical rate below which RIAF can exist.

We assume that the equipartition magnetic pressure is half the gas pressure, viz., \( B^2 / 8\pi = 0.5 \rho_i c^2 \), where \( c_s \simeq c / \sqrt{3} \times (r/r_s)^{-1/2} \) is the sound speed. This yields

\[
B \simeq 4 \times 10^4 (\dot{m} / M_9)^{1/2} (r/r_s)^{-5/4} \text{ G}. \tag{3}
\]

The numerical values in the expressions for \( n_i \) and \( B \) are in rough agreement with the results of NY95 for a viscosity parameter \( \alpha = 0.3 \), adiabatic index of 4/3, and radiative efficiency \( 1 - f \ll 1 \), where \( f \) is the advection parameter, as defined in NY95.

At radii \( r < 10^3 r_s \) the electron–ion coupling becomes weak by virtue of rapid cooling. In this region electron cooling is dominated by synchrotron emission from thermal electrons. For accretion rates near the critical rate \( \dot{m}_{\text{crit}} \), the electron temperature saturates at \( T_e \sim 10^7 \) K. At such temperatures only photons emitted by electrons at the tail of the thermal distribution have energies above the pair-production threshold. However, for much lower accretion rates, \( \dot{m} \ll \dot{m}_{\text{crit}} \), the electron temperature approaches \( 10^{10} \) K, and thermal photons can annihilate.

To estimate the annihilation rate at highly sub-critical accretion rates, we adopt the cooling functions for electron–ion and electron–electron bremsstrahlung, \( q_{ei} \) and \( q_{ee} \), from NY95. The total cooling rate per unit volume is then given to a good approximation by \( q_{\text{tot}} = q_{ei} + q_{ee} \simeq 7.5 \times 10^{-12} n_e^2 \theta_\gamma \text{ erg s}^{-1} \text{ cm}^{-3} \) at electron temperatures \( \theta_\gamma = kT_e / m_e c^2 \gtrsim 1 \). The numerical value corresponds to the choice \( \theta_\gamma = 1 \) in the logarithmic term in Equations (3.6) and (3.8) of NY95. Since the RIAF is optically thin, pair production does not affect the leptonic content of the gas, and one can safely assume \( n_e = n_i \). By employing Equation (1) one arrives at

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q_{\text{ff}} \simeq 1.8 \times 10^2 \theta_\gamma \dot{m} M_9^{-2} (r/r_s)^{-3} \text{ erg s}^{-1} \text{ cm}^{-3}. \tag{4}
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The free–free luminosity emitted by the RIAF is \( L_{\gamma \gamma} = \int q_{\gamma \gamma} d^3 r \simeq 2\pi r^3 q_{\gamma \gamma} \ln(r/r_s) \), from which we readily obtain the number density of MeV photons in the magnetosphere:

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where \( \epsilon_\gamma = 3\theta_\gamma (m_e c^2) \). The production rate of \( e^\pm \) pairs inside the magnetosphere due to \( \gamma \gamma \)-annihilation is approximately \( \sigma_{\gamma \gamma} n_e^2 c (4\pi/3) r_s^2 \). In steady state this rate is balanced by the escape rate, roughly \( 4\pi r^2 n_{\pm} c \). Equating the two rates one has

\[
n_{\pm} = \sigma_{\gamma \gamma} n_e^2 r_s / 3 \simeq 3 \times 10^{11} \dot{m} M_9^{-1} \text{ cm}^{-3}. \tag{6}
\]
The GJ density can be related to the accretion rate \( \dot{m} \) through Equation (3) with \( r = r_s \):

\[
\frac{n_{\text{GJ}}}{n_+} \simeq 6 \times 10^{12} \frac{m^{7/2}}{M_9^{3/2}} \text{cm}^{-3}.
\] (7)

Thus, we obtain

\[
n_\pm/n_{\text{GJ}} \simeq 6 \times 10^{12} \frac{m^{7/2}}{M_9^{3/2}} \frac{\rho}{r_s}.
\] (8)

As seen, below a certain accretion rate, roughly \( \dot{m} \lesssim 2 \times 10^{-4} \), injection of charges by this mechanism cannot provide complete screening of the magnetosphere. The ratio \( n_{\text{GJ}}/n_\pm \) then defines the multiplicity required to produce a force-free flow. At higher accretion rates a vacuum gap may not exist. However, as noted above, the electron temperature decreases with increasing \( \dot{m} \), eventually dropping below \( m_e c^2 \). Detailed ADAF calculations in Kerr spacetime (e.g., Mann et al. 1995; Li et al. 2009) indicate a cutoff in the emitted spectrum at around \( \theta_e \approx 1 \) even for \( \dot{m} \ll \dot{m}_{\text{crit}} \). For \( \theta_e < 1 \) the number density of MeV photons is suppressed by a factor of roughly 6 exp\((1/\theta_e)\), and \( n_\pm \propto n_+^2 \), by a factor of \( 6 \exp(1/\theta_e)^2 \). Thus, there is an uncertainty of two to three orders of magnitude in our estimate of the pair density \( n_\pm \). This large uncertainty and the strong dependence of charge injection on accretion rate motivates detailed self-consistent calculations. A complete, self-consistent treatment should also account for general relativistic effects, which are important close enough to the horizon.

The strong dependence on accretion rate suggests that a gap may form during periods of low accretion, so that emission from the gap may be intermittent.

3. TeV EMISSION AND PAIR PRODUCTION FROM CHARGES ACCELERATING IN A STARVED MAGNETOSPHERE

We consider a rapidly rotating black hole of mass \( M = 10^8 M_9 M_\odot \) embedded in an ambient radiation field. The radiation source is characterized by a luminosity \( L_d = \dot{E}_\odot \) erg s\(^{-1}\), a size \( R_d = \sqrt{\rho} \) cm, and SED (i.e., \( \nu F_\nu \)) peak energy \( E_\nu = \nu F_\nu \) eV. The energy density of the radiation field is \( \epsilon_{\text{rad}} \simeq 3 \times 10^{31} \frac{M_9}{M_\odot} \text{erg cm}^{-3} \), and the corresponding number density of photons at the peak (measured over a logarithmic energy interval) is

\[
n_\nu \simeq 1.8 \times 10^{12} \frac{M_9}{M_\odot} \text{cm}^{-3}.
\] (9)

3.1. Curvature and Inverse Compton Emission

The electric potential difference across a gap of height \( h \) generated by a maximally rotating black hole can be expressed as

\[
\Delta V = 1.7 \times 10^{21} B_4 M_9 (h/r_s)^2 \text{ Volts}.
\] (10)

Charges accelerating in the gap will quickly reach a terminal Lorentz factor at which radiative losses balance energy gain, viz., \( q \Delta V = P_t = \dot{\nu} h/c \) (Levinson 2000), where the net energy loss rate, \( P = P_\text{cur} + P_\text{IC} \), is the sum of curvature losses

\[
P_\text{cur} = \frac{2 q^2 c \gamma^4}{3 \rho^2},
\] (11)

where \( \rho \) is the curvature radius of magnetic field lines, and IC (Thomson) losses

\[
P_\text{IC} = \sigma_T c \gamma^2 u_s.
\] (12)

Equating \( \Delta V \) and \( P_\text{cur} \) and using Equations (10) and (11), one finds that curvature radiation limits the Lorentz factor of emitting electrons (positrons) to

\[
\gamma_{\text{cur}} = 5 \times 10^{10} B_4^{1/4} M_9^{3/2} (h/r_s)^{1/4} (\rho/r_s)^{1/2}.
\] (13)

If the Thomson regime applies, IC scattering of electrons on the ambient photons limits the Lorentz factor to

\[
\gamma_{\text{IC}} = \left( \frac{e \Delta V}{\sigma_T h u_s} \right)^{1/2} = 2 \times 10^9 B_4^{1/2} M_9 L_{41}^{-1/2} (h/r_s)^{1/2}.
\] (14)

Comparing Equations (13) and (14), it is seen that particle losses are dominated by IC scattering, viz., \( \gamma_{\text{IC}} < \gamma_{\text{cur}} \), if

\[
L_{41}/R^2 > 1.6 \times 10^{-3} B_4^{1/2} M_9 (h/r_s)^{1/2} (\rho/r_s)^{-1}.
\] (15)

The spectrum of curvature emission peaks at an energy

\[
\epsilon_{\text{cr,max}} = \frac{3 \hbar c^3}{2 \rho} \leq 10^{34} \frac{B_4^{1/2} M_9}{(h/r_s)^{3/4} (\rho/r_s)^{1/2}} \text{eV},
\] (16)

where \( \gamma_{\text{max}} = \min(\gamma_{\text{cur}}, \gamma_{\text{IC}}) \). The corresponding number of curvature photons emitted by a single particle is, to a good approximation,

\[
N_\gamma = P_\text{cur} (c \epsilon_{\text{cr,max}}) \geq 4 \times 10^9 B_4^{1/4} M_9^{5/2} (h/r_s)^{5/4} \times (\rho/r_s)^{-1/2} (1 + f)^{-1},
\] (17)

where \( f = P_{\text{IC}}/P_{\text{cur}} \) denotes the ratio of IC and curvature loss rates. Likewise, the maximum energy of IC photons is

\[
\epsilon_{\text{IC,max}} = m_e c^2 \frac{\gamma_{\text{max}}}{\rho} \leq 10^{34} \frac{B_4^{1/2} M_9 L_{41}^{-1/2} (h/r_s)^{3/4}}{1/2} \text{eV}.
\] (18)

The number of IC photons emitted by a single electron depends on the spectrum of the target radiation field. A crude estimate gives

\[
N_\gamma = P_{\text{IC}} (c \epsilon_{\text{IC,max}}) \geq 10^{32} B_4^{1/2} L_{41}^{1/2} (h/r_s)^{3/2} (1 + f^{-1})^{-1},
\] (19)

but the actual number may be much larger, roughly by a factor of \( m_e c^2/(\gamma_{\text{max}} h v_{\nu_0}) \) if the SED of the target radiation field peaks at a frequency \( h v_{\nu_0} \) that satisfies \( h v_{\nu_0} \ll m_e c^2 \). Inside the gap the field aligned electric field, \( E_{\parallel} \), is unscreened. This implies that the charge density on magnetic field lines must not exceed the GJ value, \( n_{\text{GJ}} = \Omega B \cos \theta/2 \pi e \). The maximum gamma-ray power that can be produced by particles accelerating in the gap, regardless of the specific gamma-ray production mechanism, is thus

\[
L_\gamma = \int n_{\text{GJ}} c (\Delta V) \pi r^2 d\theta \simeq 3 \times 10^{47} \eta B_4^2 M_9^2 (h/r_s)^2 \text{erg s}^{-1},
\] (20)

where \( \eta \) is a geometrical factor. It is worth noting that the ratio \( L_\gamma/L_{\text{BZ}} \), where \( L_{\text{BZ}} \) is the maximum Blandford-Znajek (BZ) power that can be extracted by a force-free flow, scales as \( (h/r_s)^2 \). If the pair density is well below the GJ density, then the gamma-ray power in Equation (20) is reduced by a factor of \( n_\pm/n_{\text{GJ}} \).
3.2. $\gamma\gamma$-opacity and Pair Multiplicity

The soft photon energy that corresponds to the pair-production threshold is $\epsilon_{\text{thr}} = m_e^2 c^3 / \epsilon_\gamma$, which for peak curvature photons is

$$\epsilon_{\text{thr}} = 2.5 \times 10^{-2} B_d^{3/4} M_9^{-1/2} (h/r_s)^{-3/4} (\rho/r_s)^{-1/2} \text{ eV}, \quad (21)$$

and for IC photons is

$$\epsilon_{\text{thr}} = \frac{m_e^2}{\gamma c} = 2.5 \times 10^{-4} B_d^{-1/2} M_9^{-1} L_{41}^{1/2} R^{-1} (h/r_s)^{-1/2} \text{ eV}. \quad (22)$$

Let us denote by $\zeta(\epsilon_\gamma)$ the dimensionless energy distribution of soft photons, specifically the ratio of the number density at some energy $\epsilon_\gamma$ and the number density of peak photons:

$$n_s(\epsilon_\gamma) = \zeta(\epsilon_\gamma) n_s(\epsilon_0). \quad (23)$$

Then, invoking $\sigma_{\gamma\gamma} = 0.2 \sigma_T$ and using Equation (9), the pair-production optical depth to infinity at energy $\epsilon_\gamma$ can be expressed as

$$\tau_{\gamma\gamma} = \sigma_{\gamma\gamma} n_s(\epsilon_\gamma) r_d \approx 70 \zeta(\epsilon_\gamma) L_{41} R^{-1} M_9^{-1} \zeta_0^{-1}, \quad \epsilon_\gamma \geq \epsilon_{\text{thr}}. \quad (24)$$

The optical depth across the gap is smaller by a factor $h/r_d$, where we make the reasonable assumption that $h < r_d$. Note that the optical depth depends implicitly on the gap size $h/r_s$ through the value of $\zeta(\epsilon_{\text{thr}})$ at the threshold energy.

Given $r_d$, $L_d$, and $B_d$, one can employ the above conditions in a self-consistent manner to solve for $h/r_s$ and $\gamma_{\text{max}}$. Let $d N_{\gamma}/d \epsilon_\gamma$ denote the number of photons per unit energy emitted by a single photon in the energy interval $(\epsilon_\gamma, \epsilon_\gamma + d \epsilon_\gamma)$. The gap multiplicity $\mathcal{M}$ can be defined as

$$\mathcal{M} = \int \left( 1 - e^{-(h/r_d) \tau_{\gamma\gamma}(\epsilon_\gamma)} \right) d N_{\gamma}/d \epsilon_\gamma d \epsilon_\gamma, \quad (25)$$

and the gap size is then determined from the condition $\mathcal{M} = n_{\text{GJ}}/n_\pm$, where $n_\pm$ is the density of seed charges injected into the open magnetosphere (see Equation (8)). If the latter condition cannot be satisfied it means that the gap must be fully restored, with $h/r_s \simeq 1$. An outflow may still be established above the gap if the size of the external radiation source is large enough, $r_d \gg r_s \simeq h$. In such a case the gamma-ray photons emitted by the accelerating charges in the gap will interact with target photons at larger radii, $r > h$, ensuing pair cascades there. The total multiplicity can then be computed from Equation (25) upon setting $h/r_d = 1$ in the exponent. A force-free flow will form above the starved magnetosphere provided $\mathcal{M} > n_\pm/\text{GJ}$. 

4. APPLICATION OF M87 AND Sgr A*

4.1. M87

For this source we adopt a black hole mass of $M_\bullet = 4$. The assumption that the M87 jet is extracted magnetically by a BZ mechanism imposes a lower limit on the magnetic field in the vicinity of the horizon and on the accretion rate. Various estimates (see, e.g., a summary in Li et al. 2009) yield a range of $L_j \simeq (10^{33} - 10^{44})$ erg s$^{-1}$ for the power of the M87 jet. Equating $L_j$ with the BZ power that can be extracted from a Kerr black hole having a mass $M_\bullet = 4$ and specific angular momentum $a$, $L_{\text{BZ}} \simeq \frac{\epsilon}{64} a^2 B^2 r_s^2 c \simeq 3 \times 10^{47} \epsilon a^2 m \text{ erg s}^{-1}, \quad (26)$

where $\epsilon \lesssim 1$ is an efficiency factor that depends on field topology and other details, implies an accretion rate $\dot{m} > 10^{-7} a^{-1} \text{a}^{-2}$. For a maximally rotating hole, $a = 1$, the implied accretion rate is $\sim 10^{-4}$, so that the magnetic field strength close to $r_s$ becomes $B \sim 200$ G. For an ADAF model with $L_d \sim 10 m^2 L_{\text{Edd}} (\text{NY95})$, this value for the accretion rate is consistent with limits on the bolometric luminosity measured in M87, $L_d \lesssim 10^{42}$ erg s$^{-1}$ (Owen et al. 2000). Note that the Bondi accretion rate in M87 inferred from Chandra measurements is $\dot{m}_B \sim 1.6 \times 10^{-3}$ (Di Matteo et al. 2003). The radial inflow velocity in an RIAF is about a factor $\alpha$ lower than in a Bondi flow, where $\alpha \sim 0.1$ is the viscosity parameter, so that one expects $\dot{m} \sim \alpha \dot{m}_B$ (Narayan 2002), which would be consistent with the estimate above.

The condition (8) formally yields $n_\pm/n_{\text{GJ}} = 0.12 (\dot{m}/10^{-4})^{7/2}$. However, as explained above this estimate is highly uncertain by virtue of the sensitive dependence of the injected pair density $n_\pm$ on the temperature of the accreting gas, on the geometry of the emitting region, and General Relativity effects. It seems likely, though, that a multiplicity of $10^{-2} - 10^3$ should in any case be sufficient for complete screening.

If sufficient charges cannot be provided by emission from the RIAF to ensure that $n_\pm > n_{\text{GJ}}$, then a gap will exist, where the seed charges produced by annihilation of the RIAF photons accelerate. These seed charges will emit TeV photons that can initiate pair cascades via interaction with the ambient radiation, as discussed above. However, observations constrain the $\gamma\gamma$-opacity at photon energies of a few TeV (e.g., Aharonian et al. 2006), and so pair cascades can only be initiated by photons having energies well above 10 TeV. In what follows, we employ some of the results derived in the preceding section to impose constraints on the maximum multiplicity and on the gap height.

Provided the RIAF SED is not significantly modified by the presence of an outflow, the emerging spectrum ranges from radio to X-ray energies. The radio to submillimeter regime is produced by synchrotron radiation from relativistic thermal electrons and dominated by emission from small radii. Compton up-scattering of the synchrotron photons fills in the submillimeter to hard X-ray regime, with most of the emission coming from the inner ($\lesssim 10 r_s$) region (e.g., Mann onto et al. 1997). For high $\dot{m}$, the Compton component is roughly a power law, while for low $\dot{m}$ distinct Compton peaks appear. This is related to the increase in electron temperature with decreasing accretion rates. At low $\dot{m}$, the X-ray spectrum is dominated by bremsstrahlung emission with a non-negligible contribution from large radii.

The synchrotron emission is self-absorbed, resulting in a blackbody spectrum up to frequency $\nu_{\text{sa}}$ where it becomes optically thin. In terms of ADAF parameters, one has

$$\nu_{\text{sa}}(r) \sim 9 \times 10^{10} M_9^{-1/2} (\dot{m}/10^{-3})^{1/2} (r/r_s)^{7/4} \text{ Hz}, \quad (27)$$

where $T_{\nu,\text{sa}} = T_{\nu} / 10^9$ K is the electron temperature (NY95; Mahadevan 1997). The luminosity then becomes (NY95; Mahadevan 1997)

$$\nu_{\text{sa}} L_{\nu_{\text{sa}}} \sim 9 \times 10^{35} M_9^{5/6} \left( \frac{\dot{m}}{10^{-3}} \right)^{5/6} \nu_{\text{sa}}^{7/5} T_{\nu,\text{sa}}^{21/5} \text{ erg s}^{-1}, \quad (28)$$

where $\nu_{\text{sa}} = \nu_{11} / 10^{11}$ Hz. As $T_{\nu,\text{sa}}$ varies only slowly with $r$ in the inner region of an ADAF, the highest synchrotron emission frequency comes from the innermost part ($r \sim r_s$). For M87 with $M_\bullet = 4$, $\dot{m} \sim \alpha \dot{m}_B$ and $T_{\nu,\text{sa}} \sim 5$ (cf. Mahadevan 1997),
the synchrotron peak is located in the submillimeter band at

$$v_p = v_m(r_s) \sim 10^{12} \text{ Hz} \quad (29)$$

and carries a characteristic power of

$$L_p = v_p L_{v_p} \sim 10^{41} \text{ erg s}^{-1}. \quad (30)$$

This output would still comply to the upper limit imposed by Spitzer (MIPS) observations (several arcsec resolution), indicating that the (nuclear region plus galaxy) flux in M87 at $2 \times 10^{12} \text{ Hz}$ is $\sim (3-4) \times 10^{41} \text{ erg s}^{-1}$ (Shi et al. 2007).

Compton up-scattering by a thermal distribution of relativistic electrons of small scattering depth $\tau \sim 10^{-2} (m/10^{-3}) (r_s/r)^{1/2}$, Equation (2), is characterized by a mean amplification factor per scattering of $A = 1 + \theta_B + 16\theta_B^2$. For ADAF temperatures ranges of interest, typically $2 < A \lesssim 10^2$. The frequency of the first (once-scattered) Compton peak is roughly $v_{c1} \sim \eta_L L_p \tau A$, where $\eta_L \sim 0.5$ is a geometrical factor.

Pair production implies that TeV photons of energy $\epsilon$ produced in the magnetosphere will interact most efficiently with ambient soft photons of energy $\epsilon_s \sim \epsilon (1 \text{ TeV}/\epsilon) \text{ eV}$. For moderate $A \sim 10$, the target soft photon field is dominated by the fraction of submillimeter photons scattered twice (second Compton peak) with relative peak flux $L_{c2}/L_p \sim A^2 \tau^2 \sim 10^{-2}$. Equation (24) then suggests that $\tau_{\gamma \gamma}(\epsilon) \sim 1$ up to VHE photon energies of several TeV.

If, on the other hand, $A \gg 10$, then the energy of the once-Compton-scattered synchrotron photons can approach the infrared regime, $v_r \sim 10^{12} \text{ Hz}$, yielding a non-negligible background for $\gamma \gamma$-interaction with the TeV (curvature or IC) photons. Taking $L_{c1}(\epsilon_s) \sim L_{c1} \sim 0.3$, Equation (24) with $R \sim 1$ formally implies that the pair-production optical depth for photons of $\epsilon \sim 1 \text{ TeV}$ is $\tau_{\gamma \gamma}(\epsilon) \sim 5$. However, this estimate is again somewhat uncertain by virtue of its sensitive dependence on accretion rate.

The electron temperature in the inner region of an ADAF is only weakly dependent on accretion rate and at low $m$ roughly follows $T_e \propto m^{-2} (0 < q \lesssim 0.2$ (Mannheim et al. 1997; Mahadevan 1997). The flux in the synchrotron peak, on the other hand, is highly dependent on the electron temperature and approximately varies as $L_p \propto m^{3/2} (r_s/r)^{1/2} T_E^2$ (Mahadevan 1997). For high enough temperatures $A \propto T_E^2$. This suggests that the Compton power and thus the optical depth may roughly vary with accretion rate as

$$\tau_{\gamma \gamma}(\text{TeV}) \propto L_{c1} \propto L_p A \propto m^{2.5}. \quad (31)$$

An even more sensitive dependence $\tau_{\gamma \gamma}(\epsilon) \propto m^{2.5-4}$ has been estimated based on detailed RIAF modeling of M87 (Li et al. 2009). Hence, only moderate changes in the accretion are required to make the source transparent to pair absorption up to photon energies of a few TeV as expected from IACT observations (e.g., Aharonian et al. 2006; Acciari et al. 2009). On the other hand, for energies well above 10 TeV severe absorption seems unavoidable.

The RIAF photons also provide a suitable target field for IC interactions with the energetic charged particles accelerated in the magnetosphere. As the energy loss rate for a single electron in the extreme Klein–Nishina (KN) limit is only weakly (logarithmically) dependent on the electron Lorentz factor, IC interactions (Thomson regime) with the radio-submillimeter disk photons (with the peak photons coming from $r \sim r_s$) become most constraining. IC interaction with photons of energy $\epsilon_p = h v_p$ occurs in the Thomson regime for electron Lorentz factors up to $\gamma_e \sim 10^8$. If the gap size is not too small ($h/r_s \gtrsim 10^{-2}$), electrons can get accelerated beyond this point ($P_{\text{IC}}(\gamma_e)/\epsilon < e^2 \Omega$). Taking $B_0 \simeq 0.02$ and $R \simeq 1$ for M87, IC scattering (14) would limit achievable electron Lorentz factors to $\gamma_e \simeq 10^2 L_{\gamma 21}^{-1/2} (h/r_s)_{\gamma 21}^{1/2}$, where $L_{\gamma 21} = L_{\gamma} \times 10^{41}$ erg s$^{-1}$ is the luminosity of the target field for which Compton scattering occurs in the Thomson regime ($h v_e < m_e c^2/\gamma_e$). If one takes as a first approximation

$$L_e(v) \simeq L_p \left( \frac{v_r}{v_p} \right)^{7/5}, \quad (32)$$

as inferred from Equation (28), IC interactions would limit electron Lorentz factors to

$$\gamma_e \simeq 4 \times 10^{11} L_{\rho 41}^{-5/3} (h/r_s)_{\gamma 21}^{5/3}. \quad (33)$$

Note, however, that as $v L_p \propto v^{7/2}$, with $v \propto r^{-5/4}$, the dominant contribution to $L_e$ comes from larger scales, $R_e \equiv R(v) > 1$, making the constraint even weaker, roughly by a factor of $R_{\gamma 21}^{0.1}$. Hence, if the gap is not too small, the energy loss of an electron will be dominated by curvature emission (13), limiting achievable Lorentz factors to

$$\gamma_e \simeq 3.8 \times 10^{10} \left( \frac{h}{r_s} \right)^{1/4}, \quad (34)$$

where a curvature radius $\rho = r_s$ has been invoked for the magnetic field lines.

Equation (16) then indicates that the peak energy of the curvature spectrum is limited to $\sim 50 m^{1/3} h^{-1/3}$ TeV, which is below 1 TeV for $m \lesssim 10^{-4}$. However, the maximum energy of IC photons up-scattered by the accelerated electrons exceeds $\gamma_e m_e c^2 \simeq 10^4$ TeV. Given the expected spectral index of the radio disk spectrum in M87, cf. Equation (28), the average IC power emitted by a single electron in the Thomson regime, that is, due to scattering of photons having $v < v_r = m_e c^2/\gamma_e$, is

$$P_{\text{IC}}(v) \simeq \frac{\gamma_e^2 \sigma_T L_p}{2 \pi R_d^2(v)} \left( \frac{m_e c^2}{\gamma_e^2 h v_p} \right)^{7/5} \simeq 3 \times 10^3 L_{p,41} R_t^{-2} \times (h/r_s)^{20/20} \text{ erg s}^{-1}. \quad (35)$$

These photons originate from a radius $R = R_t \gg 1$. The IC power emitted in the KN regime (for target photons at $v > v_r$) is reduced by a factor of $\ln(v_{v_p}/v)/(v_{v_p}/v)^2$:

$$P_{\text{IC}}(v) \simeq \frac{\gamma_e^2 \sigma_T L_p}{2 \pi R_d^2(v)} \ln(v/v_r) (v_{v_p}/v)^2 \simeq 6 \times 10^2 L_{p,41} R_t^{-2} \text{ erg s}^{-1} \quad (36)$$

and depends only weakly (logarithmically) on the gap height. These photons are emitted from $R = R_p \sim 1$. The total power is $P_{\text{IC}}(\gamma_e) = P_{\text{IC}}(v) + P_{\text{IC}}(v_p)$. The fraction of IC power to curvature power near the maximum energy, $I_m \equiv P_{\text{IC}}(\gamma_e)/P_{\text{cur}}(\gamma_e) \simeq 0.06 L_{p,41} R_t^{-2} (h/r_s)^{-1} \times \left[ 1 + 5(R_{p,41}/R_t)^2(h/r_s)^{20/20} \right]$, is likely to exceed a few percent.

The optical depth for the up-scattered photons at the highest energy $\epsilon_{\text{IC}} \simeq \gamma_e m_e c^2 \simeq 10^4$ TeV, cf. Equation (24) with $\xi(\epsilon_s) = (m_e c^2/\gamma_e h v_p)^{2/5}$, is likely to exceed $\tau_{\gamma \gamma}(\epsilon_{\text{IC}}) \sim 10^4$. 

\[\text{The Astrophysical Journal, 730:123 (7pp), 2011 April 1} \]
Hence, these photons will be severely attenuated before escaping the source, producing pairs. If \((h/r_s)\) is not too small, this will happen inside the gap. The arising multiplicity is of order

\[
N_{\gamma}(\text{IC}) = \frac{P_{\text{h}} c h/c}{\gamma_c m_e c^2} \simeq 10^{5.9} \frac{L_{p,41}^3 h / r_s^3}{R_p^2} \left[1 + 5(R_p/R_s)^2\right] \times \left(h/r_s\right)^{3/20} \propto m^{9/4-7\gamma}/\epsilon, \tag{37}
\]

using \(L_p \propto m^{9/4-7\gamma}\), where \(\gamma_c \propto m^{1/8}\). Typically, \(R_s \gg R_p\), so that the second term in the square parentheses can be neglected. The condition that the photon number \(N_{\gamma}(\text{IC})\) should exceed the multiplicity required to screen out the magnetosphere, \(M = n_{\text{GJ}}/n_{\pm}\), where \(n_{\pm}\) hereby refers to the seed charges, imposes a constraint on the gap height:

\[
h/r_s > \left(\frac{M}{10^3}\right)^{4/3} \frac{L_{p,41}^{-4/3} R_s^{8/3}}{r_s}. \tag{38}
\]

Note that the optical depth across the gap is \(\sim \tau_{\gamma\gamma}(\text{IC}) h/r_s = \tau_{\gamma\gamma}(\epsilon)(h/r_s)R^{-1}\) and therefore larger than unity if \(h/r_s \gtrsim 10^{-4}\).

For the above choice of parameters, Equation (20) yields a gamma-ray luminosity of \(L_{\gamma} \sim 10^{45} \eta (h/r_s)^2\) erg s\(^{-1}\), which for \(\eta = 0.1\), \((h/r_s)^2 \gtrsim 10^{-3}\) is in excess of the observed TeV luminosity. The sensitive dependence of seed charges from the RIAF emission on the accretion rate strongly suggests that moderate changes in accretion rate will induce nonlinear variations of the gap height, and conceivably its occasional disappearance. This will result in intermittencies of the gap emission and the resultant outflow.

While the curvature emission is anticipated to peak at sub-TeV energy (although any spread in the curvature radius of magnetic field lines will smear out the peak), the emitted spectrum at higher energies should depend on the details of the pair cascade process. The cascading gamma rays will be emitted from their local photosphere, at an energy that varies with radius. As demonstrated in Blandford & Levinson (1995), inhomogeneities can lead to a very broad spectrum even in the case of mono-energetic injection of electrons. Thus, detailed calculations of the cascade process are required to determine the shape of spectrum emitted from the vicinity of the gap.

### 4.2. Sgr A*\(^*\)

Estimates of the accretion rate that accommodate the extraordinary low bolometric luminosity of Sgr A*, \(L_b < 10^{-8} L_{\text{Edd}}\), are in the range \(\dot{m} \sim (0.1-5) \times 10^{-3}\), adopting a black hole mass \(M_\bullet = 4 \times 10^{-3}\) (e.g., Narayan 2002; Nayakshin 2005; Cuadra et al. 2006). VHE observations of the Sgr A* region indicate a gamma-ray spectrum that can be described by a power law extending beyond 10 TeV above threshold of 165 GeV, with an associated VHE luminosity of a few \(10^{35}\) erg s\(^{-1}\) (Aharanoin et al. 2004). The source of this emission, whether associated with the putative BH or with a different object, e.g., the PWN G359.95-0.04, is unresolved yet. From Equation (8), we obtain \(n_{\pm}/n_{\text{GJ}} \simeq 10^{-5}(\dot{m}/10^{-5})^{1/2}\).

Equation (15) implies that energy losses are dominated by curvature emission as long as \(R > 2(m/10^{-5})^{1/3} L_{36}^{1/2}\), where \(L_{36} = 1\) is roughly the bolometric luminosity observed in Sgr A*, which is dominated by emission at frequencies around the SED peak, \(v_p \sim 10^{12}\) Hz. The potential drop across the gap is limited to \(\Delta V \simeq 10^{18}(\dot{m}/10^{-5})^{1/2} V\), and from Equation (13) we deduce that the Lorentz factor of accelerating electrons is \(\gamma \sim 2 \times 10^9(\dot{m}/10^{-5})^{-1/8}\). Thus, the peak energy of curvature photons is \(\epsilon_c \sim 200(\dot{m}/10^{-5})^{3/8}\) GeV, and the maximum energy of IC scattered photons \(\epsilon_c\) is not expected to exceed \(\gamma_{\dot{m}}\epsilon_c^2 \sim 10^3\) TeV. The pair-production optical depth at this energy is \(\tau_{\gamma\gamma}(\epsilon_c) \sim 10^{5/4}\), so the IC scattered photons are likely to be converted inside the gap. Equations (3) and (19) with \(L_{\gamma} \sim 10^{-5}\) yield \(N_{\gamma} \sim 10^{-5}(\dot{m}/10^{-5})^{1/4} R^{-1}\), from which we infer that multiplicity may be high enough to allow the pair density to approach the GJ density if \(\dot{m} > 10^{-4}\) or so. For lower accretion rates we can safely assume \(h \simeq r_s\). The present estimates of \(\dot{m}\) imply that a force-free flow cannot be established by this mechanism.

The gap will, nonetheless, emit gamma rays with a total luminosity

\[
L_{\gamma} \sim 10^{35}(\dot{m}/10^{-5})^{9/2}\eta\text{ erg s}^{-1}. \tag{39}
\]

This emission is anticipated to peak in the sub-TeV band, at \(\sim 200\) GeV. Some fraction may be emitted also at much higher energies.

### 5. SUMMARY

Annihilation of MeV photons that are emitted by a radiative inefficient flow in the vicinity of a supermassive black hole leads to injection of charges into the black hole magnetosphere. The density of the injected charges depends sensitively on the accretion rate. At sufficiently large accretion rates the density may exceed the GJ density, giving rise to a complete screening of the field aligned electric field, and the formation of a force-free outflow. At lower accretion rates complete screening cannot be accomplished and a gap forms. The seed charges injected into the gap are quickly accelerated along magnetic field lines by the large potential drop generated by the rotating hole, until reaching a Lorentz factor at which energy gain is balanced by curvature and IC losses. The curvature and IC spectra peak in the TeV band. The interaction of the TeV photons thereby produced with the ambient radiation field may give rise to initiation of pair cascades, with a multiplicity that depends on the accretion rate.

The proposed scenario is applicable to variable TeV-emitting Galactic Nuclei that accrete at sufficiently low rates. Out of the nearby underluminous TeV candidate sources (Sgr A*, Cen A, M87, NGC 1399), M87 and Sgr A* are particularly interesting.\(^4\)

In the case of M87 we estimate (with a large uncertainty) that annihilation of MeV photons from the RIAF cannot provide complete screening for accretion rates that correspond to the inferred jet power, and to a fit of the observed SED by an ADAF model. However, for a fully restored gap a large multiplicity, in excess of \(10^2\), is expected owing to absorption of \(\gg\) 10 TeV photons that are produced through IC emission of the charges accelerating in the gap. The pair-production opacity decreases with increasing gamma-ray energy and becomes sufficiently small below 10 TeV to allow photons at these energies to freely escape the system. We propose that the variable TeV emission detected by H.E.S.S. was emitted from a magnetospheric gap located at the base of the VLBA jet. We note that the rapid variability is naturally accounted for by this mechanism, as even moderate changes in accretion rate will lead to nonlinear fluctuations of the gap potential and the resultant TeV emission, owing to the sensitive dependence of the density of injected

\(^4\) NGC 1399 has not yet been detected at TeV energies, while Cen A may not accrete at sufficiently low rates.
charges on accretion rate. The observed gamma-ray spectrum should depend on the spectrum of soft (scattered) photons and the pair cascade process, and can, in fact, be quite broad. Moreover, any spread in the curvature radius of magnetic field lines in the gap will lead to a spread in the maximum energy of accelerated pairs, which in turn broaden the spectrum further. Detailed calculations are needed in order to determine the shape of the emitted gamma-ray spectrum. Additional TeV photons may be emitted by the jet on larger scales, but this component should be more steady. It can be useful to trace and disentangle the spectrum of the variable component alone.

Our analysis indicates that pair production alone would lead to a very high sigma flow, which does not seem to be supported by the observations. Other processes, specifically magnetic field annihilation of a striped wind configuration by a Rayleigh–Taylor instability (Lyubarsky, 2010), might give rise to additional pair creation during the acceleration phase of the outflow. This, however, would still imply asymptotic Lorentz factor well in excess of that inferred from observations (albeit on much larger scales). Mass loading of the outflow may be accomplished through entrainment of surrounding matter, e.g., due to a non-dipolar magnetic field (McKinney & Blandford 2009), or some instability of the wall jet interface. The resulting structure may consist of a fast (high-sigma) flow, surrounded by a slower sheath (e.g., Levinson & Eichler 1993). The very inner magnetosphere may still support a vacuum gap if entrainment occurs on scales larger than several $r_g$. The details of such putative structures remain to be explored.

In the case of Sgr A* we find that for current estimates of the accretion rate the density of pairs in the magnetosphere is likely to be below the GJ value. Gamma-ray emission from the gap is expected with a total luminosity that depends sensitively on accretion rate, Equation (39), and an SED peak in the sub-TeV band.

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