Gamma-ray emission from massive stars interacting with AGN jets

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ABSTRACT

Dense populations of stars surround the nuclear regions of galaxies. In active galactic nuclei, these stars can interact with the relativistic jets launched by the supermassive black hole. In this work, we study the interaction of early-type stars with relativistic jets in active galactic nuclei. A bow-shaped double-shock structure is formed as a consequence of the interaction of the jet and the stellar wind of each early-type star. Particles can be accelerated up to relativistic energies in these shocks and emit high-energy radiation. We compute, considering different stellar densities of the galactic core, the gamma-ray emission produced by non-thermal radiative processes. This radiation may be significant in some cases, and its detection might yield valuable information on the properties of the stellar population in the galaxy nucleus, as well as on the relativistic jet. This emission is expected to be particularly relevant for nearby non-blazar sources.

Key words: galaxies: active; stars: early-type; gamma-rays: theory; radiative processes: non-thermal

1 INTRODUCTION

Active galactic nuclei (AGN) consist of a supermassive black hole (SMBH) surrounded by an accretion disc in the center of a galaxy. Sometimes these objects present radio emitting jets originated close to the SMBH (Begelman et al. 1984). These jets may be very weak or absent in radio-quiet AGN, but in radio-loud sources bipolar powerful outflows of collimated plasma are ejected from the inner regions of the accretion disc.

Radio-loud AGN produce continuum radiation along the whole electromagnetic spectrum, from radio to gamma rays. The thermal emission is radiated by matter heated during the accretion process (Shakura & Sunyaev 1973; Bisnovatyi-Kogan & Blinnikov 1972), whereas the non-thermal radiation is generated by relativistic particles accelerated in the jets (e.g. Böttcher 2007). This non-thermal emission is thought to be of synchrotron and inverse Compton (IC) origin (e.g. Ghisellini et al. 1983), although hadronic models have been also considered to explain gamma-ray sources (e.g. Mannheim 1993; Mücke & Protheroe 2001; Aharonian 2002; Reynoso et al. 2011; Romero & Reynoso 2011). In addition to continuum radiation, optical and ultra-violet emission lines are also produced in AGN. Some of these lines are broad, emitted by clumps of gas moving with velocities $v_g > 1000$ km s$^{-1}$ and located in a small region close to the SMBH, the so-called broad line region (BLR).

The presence of material surrounding the jets of AGN makes jet-medium interactions likely. For instance, the interaction of BLR clouds with AGN jets was already suggested by Blandford & Konigl (1979) as a mechanism for knot formation in the radio galaxy M87. Also, the gamma-ray production through the interaction of a cloud from the BLR with the jet was studied by Dar & Laor (1997), and more recently by Araudo et al. (2010). In the latter work, the authors showed that jet-cloud interactions may generate detectable gamma rays in non-blazar AGN, of transient nature in nearby low-luminous sources, and steady in the case of powerful objects.

In addition to clouds from the BLR, and also from the Narrow Line Region (more extended and located further away from the nucleus), stars also surround the central region of AGN. Jet-star interactions have been historically studied as a possible mechanism of jet mass-loading and deceleration in the past. In the seminal work of Komissarov (1994), the interaction of low-mass stars with jets was studied to analyze the mass transfer from the former to the lat-
ter in elliptical galaxies. Komissarov concluded that in low-
luminous jets, the interaction with stars can significantly af-
fect the jet dynamics and matter composition. In the same
direction, Hubbard & Blackman (2006) analyzed the mass
loading and truncation of the jet by interactions with stars,
also considering the case of an interposed stellar cluster.

The gamma-ray emission generated by the interaction of
massive stars with (blazar type) AGN jets has been studied
by Bednarek & Protheroe (1997). They focused on the
gamma-ray emission reprocessed by pair-Compton cascades
in the radiation field of the star, and produced by relativistic
electrons accelerated in the shocks formed by the interac-
tion of the stellar wind with the jet. Recently, Barkov et al
(2010) studied the interaction of Red Giant (RG) stars with
AGN jets, focusing on the gamma-ray emission produced
by the interaction between the tidally disrupted atmosphere
of a RG with the inner jet (see also Barkov et al. 2012;
Khangulyan et al. 2013).

In the present paper we adopt the main idea of
Bednarek & Protheroe (1997), i.e. the interaction of mas-
sove RGs with the inner jet. In this context, we consider
the radiation from jet-star interactions at different heights
(z) of the jet. We analyze the dependency on the interac-
tion region, some elliptical galaxies after a violent merger or colli-
deration processes (e.g. López-Sánchez 2010) are also expect-
ed to harbour a large number of massive stars near the active
core. Finally, star formation may take place in the external
regions of accretion discs of AGN (e.g. Hopkins & Quataert
2010), and thus a population of massive stars might exist in
the galactic core of even typical elliptical hosts.

This paper is organized as follows: in Section 2 the main
characteristics of the stellar population near the SMBH are
presented. In Sect. 3, our model of jet-star interaction is
described. In Sects. 4 and 5, the acceleration of particles and
the associated emission are studied. Then, in Sects. 6 and
7, the emission produced by the interaction of a WR and
the associated emission of X-rays is calculated, and our main
results are presented. Finally, a discussion is given in Sect. 8.

2 STELLAR POPULATIONS IN THE
NUCLEUS OF GALAXIES

The characteristics of the stellar populations surround-
ing the SMBH in AGN depend on the type of host
galaxy. Generally, in spiral galaxies the star formation rate \( \dot{M}_s \) is roughly constant, reaching values as large as
\( \sim 400 \, M_\odot \, yr^{-1} \) (Mor et al. 2012), whereas elliptical galax-
ies contain large amounts of old stars and \( \dot{M}_s \) is very low.
However, mergers between (elliptical) galaxies can lead to
renewed nuclear activity and episodes of stellar formation
(e.g. Sanders & Mirabel 1996), and accretion of matter to
the SMBH may be associated with star formation in the
galactic nuclei. In these cases \( \dot{M}_s \gtrsim 1000 \, M_\odot \, yr^{-1} \) and
the process is episodic.

The number of stars formed per mass \( (m) \), time \( (t) \) and
volume \((V \propto r^3)\) units can be expressed as
\( \psi(m,r,t) = \psi_0(m,r) \exp(-t/T) \) (Leitherer & Heckman
1993), where \( \psi_0 \equiv \psi(t = 0) \), and \( t \) and \( T \) are the age of the stellar
system and the duration of the formation process, respec-
tively. There are two limit cases: continuous formation of
stars \( (t < T) \) and starbursts \( (t \gg T) \). In the former case, \( t \) and \( T \) are the present age and the total lifetime of the host
galaxy, respectively, and being \( t \ll T, \psi \) can be considered
\( \sim \psi_0 \), and the assumption of a continuous and constant
star formation process is reasonable. In the latter case, \( t \) and \( T \) are the age and duration of the burst, respectively,
and all the stars are formed almost simultaneously, implying
\( \psi(t \gg T) \sim 0 \). In the present work we consider that
stellar formation processes take place continuously in the
nuclear region of the galaxy, and the stars are uniformly
distributed around the SMBH. The case of a jet interac-
ting with a massive star forming region will be considered
separately in a future paper.

In the present work we assume that \( \psi \) is a power-law
mass and radius distribution:

\[
\psi = K \left( \frac{m}{M_\odot} \right)^{-x} \left( \frac{r}{pc} \right)^{-y},
\]

where \( x \sim 2.3 \) for the mass range \( 0.1 \leq m/M_\odot \leq 120 \)
(Salpeter 1955; Kroupa 2001). \( y \) is a free parameter
that we fix to 1, and 2, and \( [K] = M_\odot^{-1} \, yr^{-1} \, pc^{-3} \). Mas-
sove stars are formed in giant molecular clouds with mass
\( M_* \sim 10^3 - 10^7 \, M_\odot \) and radius \( R_* \sim 10 - 200 \, pc \). Stars are
formed at a distance from the SMBH larger than the tidal

\[1\] We neglect emission produced in the shocked flows far from
the star, where there might be boosting.
\[ \gamma \text{ rays from massive stars interacting with AGN jets} \]

\[
r_t \sim 2 \left( \frac{M_{bh}}{10^7 M_\odot} \right)^{1/3} \left( \frac{M_c}{10^8 M_\odot} \right)^{-1/3} \left( \frac{R_c}{10 \text{ pc}} \right) \text{ pc.} \tag{2}
\]

with a formation rate \( \dot{M}_s = \int \int \psi_m \ dm \ \text{d}V \), i.e.,

\[
\dot{M}_s = K \int_{1 \text{ pc}}^{10^4 \text{ pc}} \left( \frac{r}{\text{ pc}} \right)^{-7/4} \pi r^2 dV \int_{0.1 M_\odot}^{120 M_\odot} \left( \frac{m}{M_\odot} \right)^{-x+1} \text{d}m.
\tag{3}
\]

To obtain \( K \), consider the empirical relation obtained by Satyapal et al. (2005):

\[
M_\odot \approx 1.0 M_{bh} c^2 \text{ is a fraction } \eta_{acc} \text{ of the Eddington luminosity, i.e., } L_{acc} = \eta_{acc} L_{Edd}, \text{ where } L_{Edd} = 1.2 \times 10^{42} (M_{bh}/10^8 M_\odot) \text{ erg s}^{-1} \text{ and } M_{bh} \text{ is the mass of the SMBH.}
\]

On Table 1 \( \dot{M}_s \) is given for \( M_{bh} = 10^7, 10^8, \text{ and } 10^9 M_\odot \), and \( \eta_{acc} = 0.01, 0.1, \text{ and } 1 \). Note that for the nine different combinations of \( M_{bh} \) and \( \eta_{acc} \), we obtain only five different values of \( \dot{M}_s \) from 0.2 to 716 M_\odot yr^{-1}. Finally, equating Eqs. (3) and (4), \( K \) results

\[
K \sim \begin{cases} 
3.22 \times 10^{-7} \eta_{acc}^{0.89} \left( \frac{M_{bh}}{10^7 M_\odot} \right)^{0.89}, & y = 1 \\
1.6 \times 10^{-8} \eta_{acc}^{0.89} \left( \frac{M_{bh}}{10^7 M_\odot} \right)^{0.89}, & y = 2.
\end{cases}
\tag{5}
\]

Once a stellar population is injected in the host galaxy, the new stars will evolve through collisions with other stars, mass loss by stellar evolution, and by stellar disruption through the loss cone (this process will enlarge \( M_{bh} \)). At the same time, stars migrate through the nuclear region forming a central cluster. Theoretical (e.g. Murphy et al. 1991; Zhao 1997) and observational (e.g. Schödel et al. 2009) studies show that stellar systems around a SMBH seem to follow a broken power-law spatial distribution \( n_s = n_s(r/r_b)^{-y} \), where \( n_s \) is the number density at the break radius \( r_b \), and \( y_1 < y_2 \) are the power-law index inside and outside \( r_b \), respectively. The presence of a SMBH produces that the most massive stars are concentrated around it and, in some cases, a stellar cusp is formed very close to the event horizon, at \( r < < r_b \), and with a slope \( \sim -0.5 \). This region is very small, but the density there is \( \sim 10^4 \) times the density predicted by \( n_s = n_s(r/r_b)^{-y} \) (Murphy et al. 1991; Zhao 1997). However, in systems with ongoing stellar formation, and low densities, relaxation timescales as tidal disruption by the SMBH and collisions between stars can be neglected. Then, stars of a given mass are accumulated in the galaxy and, at a time \( t < t_{life} \), where \( t_{life} = a(m/M_\odot)^{-b} \) is the stellar lifetime (in the main sequence), the density of stars is

\[
n_{*m} = \int_0^t \psi(t') dt' \approx \psi_0 \ t \tag{6}
\]

(Alexander 2003). For \( t > t_{life} \), stars die and the mass distribution becomes steeper than \(-2.3\), following a law \( n_{*m} \propto m^{-(2.3+b)} \). In the case of massive stars, \( t_{life} \approx (m/M_\odot)^{-1.7} \text{ and } \sim 0.1 (m/M_\odot)^{-0.8} \) Gyr, for \( 7 < m/M_\odot < 15 \) and \( 15 < m/M_\odot < 60 \), respectively (Eström et al. 2012). For \( m > 60 M_\odot \), \( t_{life} \) is almost constant and around 0.004 Gyr

\[ \text{Crowther 2012}. \]

Then, at \( t > t_{life} (8 M_\odot) \sim 0.03 \text{ Gyr, the rate of stellar formation is equal to the rate of stellar death and the system reaches the steady state for } m > 8 M_\odot. \text{ In such a case, the number density of massive stars, } n_{*m}, \text{ keeps the spatial dependence of the stellar injection rate, } \psi \propto r^{-y}, \text{ resulting}

\[
n_{*m} \sim \int_{8 M_\odot}^{120 M_\odot} \frac{0.23 \eta_{acc}^{0.89} \left( \frac{M_{bh}}{10^7 M_\odot} \right)^{0.89} \left( \frac{r}{\text{ pc}} \right)^{-y-1}}{8 M_\odot} \text{ d}m.
\]

In Figure 1 \( n_{*m} \) is plotted for the different models described in Table 1 and for the cases of \( y = 1 \) and 2. We can see from the figures that at a distance \( < 1 \) pc from the SMBH (\( \sim 10^8 R_{\text{Schw}} \) rSchw = \( 2 G M_{bh}/c^2 \) for \( M_{bh} = 10^7 M_\odot \)), the nominal density of stars is \( \sim 10^4 \) and 10 stars per pc\(^3\) for the case of \( y = 2 \) and 1, respectively. This density decreases abruptly and at a distance \( < 1 \) kpc from the center the density of massive stars would be much less than 1 star per pc\(^3\). Note that \( n_{*m} \) depends on \( \eta_{acc} M_{bh} \), and different combinations of \( M_{bh} \) and \( \eta_{acc} \) provide the same value of \( n_{*m} \).

In the next section we calculate the number of massive stars that can enter the jet, which is related to the fraction of the volume occupied by stars that is intercepted by the jet of the AGN.

3 JET-STAR INTERACTION

We are interested in the study of the interactions between massive stars and jets in AGN. In this section, we describe the main characteristics of the interaction of a massive star with a relativistic jet.

Jets of AGN are relativistic \( (v \sim c) \), with macroscopic Lorentz factors \( \Gamma \sim 5 – 10 \). The matter composition of the jets is not well known because depends on a yet incomplete jet formation theory. Two prescriptions are commonly adopted: a jet composed only by \( e^\pm \) pairs (e.g. Komissarov 1994), and a lepto-hadronic jet (e.g. Reynoso et al. 2011), i.e. \( n_p = n_e \), where \( n_p \) and \( n_e \) are the number density of protons and electrons, respectively. In such a case the jet density in the laboratory reference frame is \( \rho_j = \rho_e + \rho_p = \rho_p [1 + (m_e/m_p)] \sim \rho_p \), being \( m_e \) and \( m_p \) the rest mass of electrons and protons, respectively. Thus, we determine the jet (number) density as \( n_j = \rho_j / m_p = L_j / [(\Gamma - 1)m_p c^2 \sigma_{\gamma j}] \), where \( L_j \) and \( \sigma_j = \pi R_j^2 \) are the jet kinetic luminosity and section, respectively, and \( R_j \) its radius. According to the current taxonomy of AGN, jets from type I Faranof-Riley galaxies (FR I) are low luminous, with a kinetic luminosity \( L_j < 10^{44} \text{ erg s}^{-1} \), whereas FR II jets have \( L_j \gtrsim 10^{44} \text{ erg s}^{-1} \). The kinetic power of the jet is related with \( M_{bh} \) through the Eddington luminosity as \( L_j = \eta L_{Edd} \). In FR II sources, \( \eta \gtrsim 0.02 - 0.7 \) (Ho et al. 2008), whereas in FR I, \( \eta \lesssim 0.01 \). In the present work, we consider \( \eta < \eta_{acc} \) (see Table 1).

For the different models considered, \( L_j \) goes from \( 1.2 \times 10^{42} \) to \( 1.2 \times 10^{46} \) erg s\(^{-1} \).

Jets are probably already formed at a distance \( z_0 \sim 50 R_{\text{Schw}} \approx 5 \times 10^{-5} (M_{bh}/10^7 M_\odot) \text{ pc from the SMBH (e.g.} \text{ Junor et al. 1999). The jet expands as } R_j \sim z \tan \theta \sim \theta z, \text{ where the half opening angle } \theta \sim 1^{\circ} - 10^{\circ}. \text{ With this} \)
Table 1. Different models considered in the present work. The label assigned to each model is constructed as $M_{\text{bh}} - \eta_{\text{acccr}} - \eta_{\text{jet}}$. For instance, model M7-1-0.01 corresponds to a SMBH with mass $M_{\text{bh}} = 10^7 M_\odot \text{ yr}^{-1}$, $\eta_{\text{acccr}} = 1$, and $\eta_{\text{jet}} = 0.01$.

| $M_{\text{bh}}$ [M$_\odot$] | 20 [pc] | $L_{\text{Edd}}$ [erg s$^{-1}$] | $\eta_{\text{acccr}}$ | $M_\star$ [M$_\odot$ yr$^{-1}$] | $\eta_{\text{jet}}$ | $L_{\eta}$ [erg s$^{-1}$] | Model |
|---|---|---|---|---|---|---|---|
| $10^7$ | 5 $\times$ 10$^{-5}$ | 1.25 $\times$ 10$^{45}$ | 1.0 | 1.25 $\times$ 10$^{45}$ | 0.01 | 1.25 $\times$ 10$^{45}$ | M7-1-0.01 |
| | | | 0.01 | 1.25 $\times$ 10$^{45}$ | 0.001 | 1.25 $\times$ 10$^{45}$ | M7-1-0.001 |
| | | | 0.001 | 1.25 $\times$ 10$^{45}$ | 0.001 | 1.25 $\times$ 10$^{45}$ | M7-0.1-0.01 |
| | | | 0.001 | 1.25 $\times$ 10$^{45}$ | 0.001 | 1.25 $\times$ 10$^{45}$ | M7-0.1-0.001 |

Figure 1. Density of massive stars ($n_\star M$) and number ($N_\star j$) of early type stars inside the jet at different values of $z$ (that is equivalent to $r$), and for the case of $y = 1$ (left) and $y = 2$ (right). Cases for different combinations of $M_{\text{bh}}$ and $\eta_{\text{acccr}}$ are plotted. Other combinations not shown in the figure provide the same $n_\star M$ and $N_\star j$ plotted here. In the legend box, we did not specify the value of $\eta_{\text{jet}}$ because the plotted magnitudes are independent of this parameter.

geometry, the number of massive stars contained inside the jet volume $V_j$ is $N_\star j(z) = \int_{V_j} n_\star M(z')dV_j$, where $dV_j = \pi R_j^2 dz' (z$ is the coordinate along the jet). This yields:

$$N_\star j(z) \sim \begin{cases} 2.89 \eta_{\text{acccr}}^{0.89} \left( \frac{M_{\text{bh}}}{10^7 M_\odot} \right)^{0.89} \left( \frac{z}{R_j^2} - 1 \right), & y = 1 \\ 1.43 \times 10^{3} \eta_{\text{acccr}}^{0.89} \left( \frac{M_{\text{bh}}}{10^7 M_\odot} \right)^{0.89} \left[ \left( \frac{z}{R_j^2} \right) - 1 \right], & y = 2 \end{cases}$$

At $z \geq z_1 \sim 1.6\eta_{\text{acccr}}^{0.89}(M_{\text{bh}}/10^7 M_\odot)^{-0.89}$ and $n_{\text{acccr}}^{0.89}(M_{\text{bh}}/10^7 M_\odot)^{-0.89}$ pc, for the case of $y = 1$ and $y = 2$, respectively, there is at least one massive star inside the jet at every time (see Fig. 1). For $z$-values such that $N_\star j < 1$, then $N_\star j$ is the fraction of time during which there is a star within the jet.

The permanence of stars inside the jet is determined by the jet crossing time $t_j \sim 2R_j/v_\star \sim 7 \times 10^5 (z/pc)^{3/2} (M_{\text{bh}}/10^7 M_\odot)^{-1/2}$ yr, where $v_\star = (2G M_{\text{bh}}/z)^{1/2} \sim 3 \times 10^7 [(M_{\text{bh}}/10^7 M_\odot) (z/pc)]^{-1/2}$ cm s$^{-1}$ is the speed at which stars are moving around the SMBH. To analyze the interaction of stars with the jets, we need to know the structure of the shocks formed as a consequence of the collision of the jet plasma with the stellar wind. The double bow shock formed around the star (see Komissarov 1994 for a detailed study of the bow-shock structure and stability) depends not only on the jet properties, but also on the stellar ones, in particular the stellar wind mass-loss rate and velocity. Main-sequence massive (OB) stars have typically mass-loss rates $M_\star \sim 10^{-7} - 10^{-5}$ $M_\odot$ yr$^{-1}$. This mass-loss is radiatively driven, forming supersonic winds that reach terminal velocities $v_\infty \sim 3000$ km s$^{-1}$ in the fastest cases (e.g. Lamers & Cassinelli 1999). The luminosities and surface temperatures of OB stars are $L_\star \sim$.
Considering that the wind velocity is described by a $\beta$-law, i.e. $v_w = v_\infty(1 - R/R_*)^{\beta}$, where $\beta \sim 0.8 - 1$, at distances $R \gtrsim 2 R_*$, the approximation $v_w \sim v_\infty$ is reasonable.
A. T. Araudo et al.

A. T. Araudo et al. (2009, and references therein, for a derivation and ˙ΓM, are not so different: Lbs/Lw = 4.4(107 M⊙)0.06, y = 1

Since we neglect flow reacceleration downstream the bow shock, or shading of shocks by other shocks further up-stream, when the energy rate crossing all the shocks reaches ~ L1 (i.e. τ = 1), the jet is completely stopped. When τ > 1, the approximation of a constant L1 is not valid any more. However, this occurs at z > 1 kpc on our models.

With the reduction of L1 by jet-star interactions, the jet velocity will decrease. For a cold jet Lj = Mj(Γ − 1)c2 and considering Mj constant, the Lorentz factor results Γ = Γ0 exp(−τ) + 1. However, the assumption of constant Mj is not strictly correct. The entrainment of cold material from the stellar wind will also contribute to the deceleration of the jet bulk motion. In the surface discontinuity a mixing layer will develop, and the shocked jet and wind matter can mix. This mixing is produced by turbulent motions in the bow shock tail, likely triggered by Rayleigh-Taylor (RT) and Kelvin-Helmholtz (KH) instabilities. Under effective mixing, Mj(z) = Mj0 + Mj(z), where Mj0 is the initial rate of jet mass and Mj(z) ∼ Nj(z)Mw. This effect has been analyzed by Komissarov (1994) for the case of low-mass stars (typical Mw ∼ 10−12 M⊙ yr−1) interacting with jets, concluding that mixing by KH instabilities is an important mechanism of mass loading in FR I galaxies. In the next subsection we show, through a simple analysis of timescales, that KH instabilities are also important in the case of massive stars interacting with jets.

3.2 Dynamical timescales

We are interested in the bow shocks generated around the stars as places for acceleration of particles and production of non-thermal emission. For this reason, even when we will not study the dynamics of these bow shocks, we will estimate the time during which stars can be inside the jet as obstacles, and the evolution and interplay of the shocked flows.

The time required by the star and its wind to penetrate into the jet is tbs ~ 2Rsp/vw ~ 5.6 × 102(z/z0)3/2 s. In addition to tbs and t1, there are also hydrodynamical instabilities produced by the jet interaction that affect the shocked flows, triggering their disruption and mixing. The timescale for full development of the two bow-shock structures is roughly tba ~ Rsp/csw, where csw is the sound speed in the wind shock, csw ~ vw. This is also the timescale on which RT and KH instabilities will lead to irregularities in the contact discontinuity of size ~ Rsp (see, e.g., Araudo et al. 2009, and references therein, for a derivation of these timescales); RT mainly acting in the region around the apex of the contact discontinuity, and KH in the outflowing tail, further downstream. For effective disruption of the two shocked flows, and their acceleration by the jet thrust and eventual mixing, a time of order of few times tbs is needed, which yields a length for the mixing tail of about few times tbsv1/vj ~ Rspχ1/2, where χ ≡ vj/vw ≈ √ρw/Γj. If the ratio Rj/Rsp is of the order of or larger than χ, then jet dilution with z will not have a relevant impact on the process. Otherwise, jet dilution will likely weaken the instability growth on the largest tail scales, slowing down mixing. Effective mixing also requires that tbs < t1, since otherwise the interaction structure will not fully develop. Given the values of Rsp, Rj, vj and vw considered in this work, the mixing conditions seem to be fulfilled, and larger Mh0-values (implying larger vw) should not have a significant impact. For simplicity, we have kept the reasoning at a basic level. For a more accurate and detailed description of tail disruption within relativistic jets, we refer to Blandford & Königl (1979) and Komissarov (1994).

4 NON-THERMAL PARTICLES

In addition to the dynamical processes described above, non-thermal particles can be generated in jet-star interactions. In the bow shocks, particles can be accelerated up to relativistic energies through a Fermi-like type I acceleration mechanism. The size of the jet and wind shocked regions, Dj and Dw, respectively, are determined considering the conservation of the rate of the number particle density. Using relativistic and non-relativistic Rankine-Hugoniot relations we obtain that Dj ∼ Dw ∼ 0.3Rsp.

Although the jet kinetic luminosity is much larger than the wind luminosity (Lw = Mwvw2/2) at the location of Rsp,

\[ \frac{Lj}{Lw} = 5 \times 10^{45} \left( \frac{Lj}{10^{42} \text{ erg s}^{-1}} \right) \left( \frac{Mw}{10^{-6} \text{ M}_\odot \text{ yr}^{-1}} \right)^{-1} \times \left( \frac{vw}{2000 \text{ km s}^{-1}} \right)^{-2}, \]

the available luminosity in the jet and wind bow shocks, Ljbs and Lwbs, respectively, are not so different: Ljbs/Lwbs ∼ 75 (vw/2000 km s−1)−1.

\[ \text{We have considered the Rankine-Hugoniot relations obtained for the case of a plane shock.} \]
A fraction \( \eta_{\text{jet}} \) of these luminosities is transferred to particles accelerated in each shock, implying a non-thermal luminosity in the jet \( L_{\text{nj}} = \eta_{\text{jet}} L_{\text{bs}} \), and in the wind \( L_{\text{nw}} = \eta_{\text{jet}} L_{\text{wbs}} \). The fraction \( \eta_{\text{jet}} \) is a free parameter of the present model. We assume that the populations of accelerated electrons and protons have the same luminosity, and we fix \( \eta_{\text{jet}} = 0.1 \) both in the jet and in the wind bow shocks. We note that the radiation luminosity scales simply as \( \propto \eta_{\text{jet}} \).

Relativistic particles are assumed to be injected in the downstream region of the bow shocks following a power-law energy distribution: \( Q_{e,p} \equiv K_{e,p} E_{e,p}^{-\gamma} \exp(-E_{e,p}/E_{e,p}^{\text{max}}) \), where \( E_{e,p}^{\text{max}} \) is the maximum energy achieved by particles, and \( e \) and \( p \) stands for electrons and protons, respectively. A power-law index \( \gamma \simeq -2 \) is usual for Fermi I-type acceleration mechanisms, and the normalization constants \( K_{e,p} \) are determined through \( \eta_{\text{jet}} = \int Q_{e,p} E_{e,p} dE_{e,p} \).

As a consequence of radiative and escape losses, the injected particles evolve until they reach the steady state, with characteristic timescales \( t_{\text{adv},j} \sim 4 R_{\text{bs}}/c \) and \( t_{\text{adv},w} \sim 4 R_{\text{bs}}/v_{\infty} \), i.e. the advection timescale in the downstream regions of the jet and the wind bow shocks, respectively. In this work we consider that the emitting regions are uniform, i.e. we adopt a one-zone model for the accelerator/emitter. Under this condition, we solve the following equation to derive the energy distribution of relativistic electrons and protons, \( N_{e,p} \):\(^{(196)}\)

\[
N_{e,p} \left( \frac{t_{\text{esc}}}{t_{\text{diss}}} \right) \left( \frac{E_{e,p}}{E_{e,p}^{\text{max}}} \right) = Q_{e,p},
\]

where \( t_{\text{esc}} = \min(t_{\text{diss}}, t_{\text{adv}}) \). The diffusion timescale is \( t_{\text{diss}} \sim D_{e,p} q B_{\text{bs},wbs}/(E_{e,p} c) \) in the Bohm regime, where \( B_{\text{bs}} \) and \( B_{\text{wbs}} \) are the magnetic fields in the jet and the stellar wind bow-shock regions, respectively, and \( q \) is the electron charge. In addition to diffusion, particles suffer different relevant radiative losses \( E_{e,p} \), synchrotron and stellar photon IC upscattering for electrons, and proton-proton (Kelner et al. 2006). All mentioned losses balance the energy gain from acceleration, \( E_{e,p} \), when the steady state is achieved.

4.1 Particle acceleration and losses in the jet shock

The fraction of the jet section that is intercepted by the stellar bow shock, \( \eta_{j} = \sigma_{\text{sp}}/\sigma_{j} \), is \( \propto L_{j}^{-1} \). Therefore, \( L_{\text{bs}} = \eta_{j} L_{j} \) results to be independent of \( L_{j} \) and \( z \):

\[
L_{\text{bs}} = \left( \frac{R_{\text{bs}}}{R_{j}} \right)^{2} L_{j} \sim 10^{38} \left( \frac{M_{\odot}}{10^{-6} M_{\odot} \text{ yr}^{-1}} \right) \left( \frac{v_{\infty}}{2000 \text{ km s}^{-1}} \right) \text{ erg s}^{-1}.
\]

Note however that for rare cases of stars interacting at \( z < z_{j} \), \( R_{\text{sp}} = R_{j} \) and \( L_{\text{bs}} = \propto L_{j} \) \( z^{-2} \). The jet bow shock has a velocity \( \sim v_{j} \), and particles are accelerated at this relativistic shock with a rate assumed to be \( E_{e,p} ^{\text{acc}} \sim 0.01 q B_{\text{bs}} c \). We adopt a modest acceleration efficiency, although for relativistic shocks, values as high as \( E_{e,p} ^{\text{acc}} \sim \sim 0.1 q B_{\text{bs}} c \) have been derived (e.g. Achterberg et al. 2001).

Theoretical studies on jet acceleration (e.g. Komissarov et al. 2007) suggest that near the base of the outflow, the kinetic energy density of the jet, \( U_{\text{kin}} = L_{j}/(\sigma_{j} v_{j}) \), is smaller than the magnetic energy density \( U_{\text{mag}} = B_{j}^{2}/8\pi \), where \( B_{j} \) is the jet magnetic field. However, at \( z \gtrsim 10^{-3} M_{\text{bs}}/10^{7} M_{\odot} (\theta/5^\circ)^{-1} \text{ pc} \), magnetic forces have already accelerated the flow and \( U_{\text{kin}} \) is likely to be dominant. Given that we are interested in the jet properties at \( z \gtrsim 1 \text{ pc} \), we estimate \( B_{j} \) assuming that \( U_{\text{mag}} \sim 0.3 U_{\text{kin}} \), with \( \eta_{\text{mag}} = 0.3 \) (see Fig. 8 of Komissarov et al. 2007) for the case of conical jets). The corresponding magnetic field is:

\[
B_{j} \sim 0.34 \left( \frac{\eta_{\text{mag}}}{0.3} \right)^{1/2} \left( \frac{L_{j}}{10^{42} \text{ erg s}^{-1}} \right)^{1/2} \left( \frac{q}{5^\circ} \right)^{-1} \left( \frac{z}{\text{pc}} \right)^{-1} \text{ G}.
\]

Then, assuming that in the bow shock downstream region is amplified by the compression of the flow, \( B_{\text{bs}} \) results to \( 4 B_{j} \sim 4.1 \left( \eta_{\text{mag}}/0.3 \right) (L_{j}/10^{42} \text{ erg s}^{-1})^{1/2} (z/\text{pc})^{-1} \text{ G} \).

The most important radiative losses of relativistic electrons in the jet bow-shock region are synchrotron and IC scattering of photons from the star. At \( R_{\text{sp}} \), the energy density of photons is \( U_{\gamma} \sim L_{\gamma}/(4\pi R_{\text{sp}} c) \sim (L_{j}/10^{42} \text{ erg s}^{-1}) (z/\text{pc})^{-2} \text{ erg cm}^{-3} \). Considering that these photons follow a thermal distribution with a maximum at an energy \( E_{\gamma} \sim 3 K_{B} T_{\gamma} \sim 10(T_{\gamma}/3 \times 10^{4} \text{ K}) (K_{B} = 1.4 \times 10^{-16} \text{ erg K}^{-1} \text{ is the Boltzmann constant}), at \( |E_{\gamma}| > (mc^{2})^{2}/E_{\gamma} \sim 50 \text{ GeV} \), the IC interaction occurs in the Klein-Nishina (KN) regime. Photons from the accretion disc are a less important target for IC interactions compared with photons from the star, as seen from the large value of the ratio \( U_{\gamma}/U_{\text{d}} \sim 10^{2} \), for the wind parameters adopted here and adopting a disc luminosity \( \sim L_{j} \). Electrons can also radiate through relativistic Bremsstrahlung in interactions with the shocked jet matter. Nevertheless, densities are so low that relativistic Bremsstrahlung losses are quite inefficient when compared with escape, synchrotron or IC scattering.

The maximum energy achieved by electrons in the jet shock is determined by synchrotron losses resulting

\[
E_{\text{max}} \sim 20.3 \left( \frac{\eta_{\text{mag}}}{0.03} \right)^{-1/4} \left( \frac{z}{\text{pc}} \right)^{1/2} \left( \frac{L_{j}}{10^{42} \text{ erg s}^{-1}} \right)^{-1/4} \text{ TeV}.
\]

In Fig. \( \text{H} \) \( E_{\text{max}} \) is plotted for different values of \( L_{j} \). Taking into account the escape, synchrotron, and IC losses described above, we solve Eq. \( \text{H} \) obtaining the energy distribution \( N_{e} \) of relativistic electrons shown in Fig. \( \text{H} \) (left). The synchrotron and IC cooling dominate the high-energy part of the electron energy distribution, and at low energies advective losses are dominant. This appears as a steepening in \( N_{e} \) as \( \propto E^{-2.1} \) to \( \propto E^{-3.1} \).

The maximum energy of protons accelerated in the jet shock is determined by advection losses, giving \( E_{\text{max}} \sim 2 \times 10^{9} (R_{\text{bs}}/0.03)^{-1/2} \text{ TeV} \). These relativistic protons escape from the jet bow-shock region advected by shocked matter, without producing significant levels of radiation. For this reason we do not take into account hadronic emission from the jet shocked region. Given that the proton energy is below the photomeson production threshold with stellar photons as targets, this process can also be neglected.

\( \gamma \) rays from massive stars interacting with AGN jets
4.2 Particle acceleration and losses in the wind shock

Assuming that the whole wind is shocked, the shock luminosity would be:

\[ L_{\text{whs}} \approx 1.3 \times 10^{36} \left( \frac{M_w}{10^{-6} \, M_\odot \, \text{yr}^{-1}} \right) \left( \frac{v_w}{2000 \, \text{km s}^{-1}} \right)^2 \text{erg s}^{-1}. \]  

(19)

Being this shock non-relativistic, with velocity \( v_w \), particles are accelerated with a rate \( E_{\text{ac},p} = (1/2\pi) q(v_w/c)^2 B_w v \) (e.g. Drury 1983).

The magnetic field of the wind, \( B_w \), roughly has a dipolar structure close to the star surface, and radial and toroidal components dominate farther out (Usov & Melrose 1992). For simplicity, we will adopt here \( B_{\text{whs}} \approx B_w \). Fixing \( B_w = 10 \, G \), \( B_{\text{whs}} \) results \( \approx 0.1 B_{\text{whs}} \) at \( z > z_* \), and synchrotron cooling will be more efficient in the jet than in the wind shocked region. On the other hand, given that the size and radiation field values are similar, the IC cooling timescale in the wind shocked region is similar to the one in the jet. The main difference is in the advection timescale and the maximum energy, given the much lower shock velocity. The lower advection speed implies that the electron energy distribution steepens at lower energies, implying a high radiation efficiency. The maximum energy of electrons accelerated in the wind is determined by IC and diffusion losses, providing the values of \( E_{\text{emax}} \) plotted in Fig. 5( right), and it is similar to the distribution of electrons accelerated in the jet, i.e. at low values of \( z \) \( E_{\text{emax}} \propto E_{\gamma}^{-2.1} \) as a consequence of IC and synchrotron losses, with a hardening beyond \( \sim 10 \, \text{GeV} \). At larger heights, \( E_{\gamma} \propto E_{\text{emax}}^{-2.4} \) all the way up to \( E_{\text{emax}} \) as a consequence of advection escape losses.

Regarding protons, the large wind particle densities imply that the proton-proton cooling channel is more efficient than in the jet shocked region, but still it is a minor channel of gamma-ray production compared with IC for the same \( e \) and \( p \) energetics. The proton energy distribution is dominated by advection losses, which are independent of energy, and therefore it keeps the injection slope, i.e. \( N_p \propto E_{\gamma}^{-2} \). The maximum energy of protons is constrained by diffusion losses, giving \( E_{\text{emax}}^p = 0.2(B_w/10 \, G)(v_w/2000 \, \text{km s}^{-1}) \, \text{TeV}. \)

4 We cannot provide an analytical expression for \( E_{\text{emax}} ^{\gamma} \) in the wind because in the range where it is constrained by IC scattering in the KN regime, the calculation was done through the Runge-Kutta numerical method.
5 NON-THERMAL EMISSION

Once \( N_z \) in the jet and wind shocked regions is computed, we calculate the spectral energy distribution (SED) of the non-thermal radiation, synchrotron and IC scattering (in Th and KN regimes) in the jet and the wind shocked regions, using the standard formulae (e.g. Blumenthal & Gould 1970). The energy budget for the emission produced in the bow shock regions are \( \eta_{\text{d}} L_{\text{obs}} \), and \( \eta_{\text{d}} L_{\text{wbs}} \), which would be an upper limit for the emission luminosity produced both in the jet and in the wind, respectively.

An important characteristic of the scenario studied in this paper is that the emitter is fixed to the star, and being the star moving at a non-relativistic velocity, the emission produced in the bow shock regions is not amplified by Doppler boosting.

At radio wavelengths, the synchrotron self-absorption effect has been taken into account, although it is only relevant for interactions very close to the jet base. At gamma-ray energies, photon-photon absorption due to the presence of the stellar radiation field can be relevant at certain z (e.g. Bednarek & Protheroe 1997), but the internal absorption due to synchrotron radiation is negligible. Given the typical stellar photon energy \( E_0 \sim 10 \text{ eV} \), gamma rays beyond \( \sim 30 \text{ GeV} \) can be affected by photon-photon absorption. However, this process is only important at small z, where \( R_{\text{sp}} \) is also small. At \( z > 1 \) pc SEDs shown in Fig. 6 are not strongly absorbed. Another effect that should be taken into account at energies beyond 100 GeV is absorption in the extragalactic background light via pair creation (important only for sources located well beyond 100 Mpc).

The leptonic emission is indistinguishable if \( R_{\text{sp}} \) is the same, regardless the \( z \) of interaction and \( L_j \). However, more powerful jets have a transition from radiation to advection dominated interactions at higher \( z \)-values, which enhances their non-thermal luminosity. Synchrotron and IC losses are proportional to magnetic (energy) and radiation densities, and thus are \( \propto z^{-2} \). The increase of the time during which particles remain in the emitter, \( \propto z \), and the growth of the number of stars within a jet slice, \( \propto z^{0.25} \), are not enough to balance the loss in radiation efficiency beyond the \( z \) at which radiation cooling is not dominant (at any particle energy). This implies that there is more emission generated at relatively small \( z \)-values. To illustrate the changes in the SED with \( z \), we present in Fig. 6 the synchrotron and IC emission produced by the interaction of only one star with the jet at \( z = 1, 10, 100 \), and \( 10^4 \) pc, adopting the parameters listed on Table 3 (right panel) for the different models presented on Table 1. A detailed description of Fig. 3 is given in Sections 5.1 and 5.2. In addition to that, we calculate the bolometric luminosities achieved by synchrotron and IC emission in the jet -\( L_{\text{IC,pen}} \)- and in the wind -\( L_{\text{IC,adv}} \)- by the interaction of only one star at different \( z \) from 1 pc to 1 kpc. In Fig. 6, \( L_{\text{IC,pen}} \) and \( L_{\text{IC,adv}} \) (maroon-solid lines) are shown. A detailed description of this figure is given in Sections 5.1 and 5.2 and also in Sec. 7.

5.1 Leptonic emission from the jet shock

The synchrotron and IC emission from the jet bow shock is presented in Fig. 6 (left panel). As mentioned, both synchrotron and IC are more efficient in the inner jet regions, emission at lower energies getting less efficient (due to advection) at higher \( z \)-values. This effect is clearly seen in Figs. 6 and 7 (both in the left panel).

In Fig. 6 we see different spectral features in the cases of \( L_j = 1.2 \times 10^{42} \) (top) and \( 1.2 \times 10^{46} \) erg s\(^{-1}\) (bottom). In the former case, the break energy in the photon spectrum is higher than in the latter case. This is very clear in the synchrotron emission, where the break energy in the case of \( L_j = 1.2 \times 10^{42} \) erg s\(^{-1}\) is about \( 10^8 \) times larger than in the case of \( L_j = 1.2 \times 10^{46} \) erg s\(^{-1}\). (Compare Figs. 6 and 7). Another clear difference is the break produced by synchrotron self-absorption, being the source optically thick at lower energies in the case of \( L_j = 1.2 \times 10^{42} \) erg s\(^{-1}\) than in the case of \( L_j = 1.2 \times 10^{46} \) erg s\(^{-1}\). Photon-photon absorption in the IC spectrum is not relevant in any case.

The total bolometric luminosity produced in the jet, \( L_j = L_{\text{IC,pen}} + L_{\text{IC,adv}} \), where \( L_{\text{IC,pen}} \) and \( L_{\text{IC,adv}} \) are the bolometric luminosities of IC and synchrotron radiation in the jet, respectively, is plotted on the left panel of Fig. 6 (maroon-solid line). Note that as \( z \geq 1 \) pc, where\( R_{\text{sp}} \propto z \), \( L_{\text{IC,adv}} \sim 10^{37} \) erg s\(^{-1}\) is constant on \( z \) as is shown in Fig. 7 with a black-solid line.

5.2 Leptonic emission from the wind shock

The synchrotron and IC emission from the wind bow shock is presented in Fig. 6 (right panel), also for the cases of only one star interacting with a jet of \( L_j = 1.2 \times 10^{42} \) (top) and \( 1.2 \times 10^{46} \) erg s\(^{-1}\) (bottom), at \( z = 1, 10, 100 \), and \( 10^4 \) pc. The SED shows lower maximum energies and lower achieved emission levels than those of the shocked jet region. We can see from the figure that the synchrotron emission produced in the wind is very faint, with an specific luminosity about five order of magnitude lower than the IC emission.

The total bolometric luminosity produced in the wind, \( L_w = L_{\text{IC,pen}} + L_{\text{IC,adv}} \), is plotted on the right panel of Fig. 6 (maroon-solid line). Note that at \( z \geq 1 \) pc, \( L_w \propto z^{-2} \). Finally, note that as a consequence of \( L_{\text{IC,pen}} / L_{\text{IC,adv}} \sim 100 \), because \( \nu_{\text{cut}} / \nu_{\text{IC}} \sim 100 \), the fraction of the available non-thermal luminosity that is radiated in the wind is larger than that of the jet emission, i.e. \( L_w / L_{\text{IC,adv}} > L_j / L_{\text{IC,adv}} \).

Although \( n_w \) is larger than in the shocked jet region, the production of gamma rays by proton-proton interactions of wind accelerated protons and shocked matter is negligible when compared with emission from IC scattering. For this reason, we do not compute the luminosity produced by this emission channel. Besides that, the synchrotron and IC emission from e\(^{+}\) secondaries of these proton-proton interactions will be much smaller than that from primary electrons.

6 FLARING EMISSION FROM A WOLF-RAYET STAR

Wolf-Rayet stars evolve from OB-type stars. Typically, WR stars have masses \( \sim 10 - 25 \text{ M}_\odot \), and strong mass-loss rates, \( \sim 10^{-4} \text{ M}_\odot \text{ yr}^{-1} \). They are very luminous, \( L_{\text{WR}} \sim 10^{49} \) erg s\(^{-1}\), reaching photospheric radius as large as \( 10^2 \text{ R}_\odot \) in the most powerful cases (Crowther 2007). Since WR stars are scarce, it is not expected to find large populations of WR stars in the inner region of AGN, we will consider here the situation of a single WR star interacting...
with the AGN jet. The winds of WR stars are so powerful that can balance the ram pressure of a jet with \( L_j = 1.2 \times 10^{42} \text{ erg s}^{-1} \) at any \( z \), since \( z \sim 0.74 \left( \frac{L_j}{10^{42}} \text{ erg s}^{-1} \right)^{1/2} z_0 \) for the properties of the WR star listed on Table 3.

In order to compare the spectrum produced by a WR star and by standard massive (OB) stars as was shown in Section 5, we assume that the WR penetrates the jet at \( z = 1 \text{ pc} \) in the case with \( L_j = 1.2 \times 10^{42}, 1.2 \times 10^{44}, \) and \( 1.2 \times 10^{46} \text{ erg s}^{-1} \). Being \( M_{WR}/M_* = 100 \), the stagnation point of the WR wind is located at \( R_{sp,WR} \sim 10R_{sp} \). Thus, the available luminosity to accelerate particles in the shocks produced by the interaction of the WR is \( \sim 100 \) times larger than in the case of an OB star. In Fig. 8 we show the synchrotron and IC emission produced in the jet and in the wind. Note that the IC emission from the wind reaches similar levels to the IC emission from the jet, on the contrary to the case of an OB star, where the IC jet emission in the case of \( 1.2 \times 10^{46} \text{ erg s}^{-1} \) is \( \sim 100 \) times smaller in the wind than in the jet. This is a consequence of the different energy breaks in the electrons energy distribution. Comparing the curves that correspond to \( z = 1 \text{ pc} \) in Fig. 8 with Fig. 8 we can appreciate that the shape of synchrotron and IC spectrum in the jet is different for the case of an OB star and a WR, where in the former case the break energy produced by the advection escape to the radiation dominated regime is at higher energies than in the latter. Finally, the emission level produced by a WR (both in the wind and in the jet) is larger than the one produced by an OB star (both interacting with the jet at the same \( z \)).

The radiation produced by a WR interacting occasionally with a jet will be transient with a timescale similar to the jet crossing time, unlike the steady emission produced by a population of stars, described in the next section. We remark that, if the star diffusion time were short enough to allow a massive star to reach the vicinity of the SMBH in

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**Table 3. Parameters of the WR star considered in this work.**

| Description                  | Value                      |
|------------------------------|----------------------------|
| Mass loss rate               | \( M_{WR} = 10^{-4} M_\odot \text{ yr}^{-1} \) |
| Wind terminal velocity       | \( v_{WR} = 3000 \text{ km s}^{-1} \) |
| Luminosity                   | \( L_{WR} = 10^{49} \text{ erg s}^{-1} \) |
| Surface temperature          | \( T_{WR} = 3 \times 10^4 \text{ K} \) |
Figure 7. Bolometric luminosities (maroon-solid lines) in the jet (left) and in the wind (right), produced by the interaction of only one star along the whole jet: from 1 pc to 1 kpc. The thickness of maroon lines is increased from low to large values of $L_j$, as is indicated in the right panel. The bolometric luminosity of many stars up to a certain $z$ is also presented. Cases with $M_{bh} = 10^7$ (top), $10^8$ (middle), and $10^9 M_\odot$ (bottom) are shown. In each plot, the results of the two spatial distribution models ($y = 1$ -green lines- and 2 -red lines) are presented. The thickness of green and red lines is increased from low to large values of $\eta_{accr}$. The black line indicates the value of $L_{jbs}$ (left), and $L_{wbs}$ (right).
achieve levels of z > z1, there are more than one star every time into the jet, and the non-thermal luminosity produced by all the stars into the jet at different z is also plotted in Fig. 4. In each panel we show the bolometric luminosity produced by the different stellar populations considered in the present study. Note that on the one hand, in the most powerful case (M9-1-0.1) the total bolometric luminosity produced in the jet and in the wind is \( \sim 5 \times 10^{41} \) and \( \sim 10^{39} \text{ erg s}^{-1} \) (y = 2), respectively. On the other hand, in cases with low density of massive stars, the luminosity produced by the cumulative effect of all stars into the jet can be lower than the luminosity produced by only one star interacting with the jet close to z0 if the star formed at \( z \gtrsim r_t \) and migrated close to the jet base.

Considered the \( M_\ast - M_{bh} \) relation given by Satyapal et al. (2003), the density of massive stars results \( \propto (\eta_{acc} M_{bh})^{0.69} \). Thus, sources with \( M_{bh} = 10^8 - 10^9 M_\odot \) and \( \eta_{acc} \sim 1 - 0.1 \) are likely to be detected at gamma rays by Fermi with a deep enough (pointed) exposure or after some years of observation in the survey mode. In the case of stellar populations around a SMBH with \( M_{bh} = 10^7 M_\odot \), the gamma ray emission produced cannot be detected by Fermi in any case, and the same occurs for \( M_{bh} = 10^8 M_\odot \) and \( \eta_{acc} \sim 0.01 \) (under the assumed \( \eta_B \) and \( \eta_{acc} \)). The most interesting case is that of a high accretion rate \( \eta_{acc} \sim 1 \) and \( M_{bh} = 10^9 M_\odot \) yr\(^{-1} \), whose emission can be detected in the case of luminous \( L_j \sim 10^{46} \text{ erg s}^{-1} \) and close sources (such as M87). Less luminous sources may also be detected in the near future by the Cherenkov Telescope Array (CTA).

In the case of a population of massive stars (continuously) interacting with the jet, the produced emission will be steady and produced in a large part of the jet volume, from \( z_1 \) to \( z_2 \), on scales of \( \sim \text{kpc} \).

7 STEADY EMISSION FROM A POPULATION OF MASSIVE STARS

In order to study the emission produced by many massive stars, we assume within the jet a stellar population as the one described in Sect. 2. As shown in Sect. 6 the emission produced at small values of \( z \) is higher than the emission produced at larger \( z \), as a consequence of the dilution of the target fields with \( z \). This effect is balanced by the fact that, at \( z > z_1 \), the number of stars interacting with the jet is \( > 1 \) and the emission produced by all of them increase \( \propto z^2 \) and \( \propto z \), for the cases with \( y = 1 \) and \( 2 \), respectively. We calculate the emission produced by each of the \( N_j \) stars at a certain \( z \), and then integrate along \( z \) all the contributions, obtaining the SEDs shown in Fig. 3 for different values of \( M_{bh} \) and \( L_j \). Note that the features of these SEDs are similar to the SED produced by only one star located at a relatively large value of \( z \) (see Fig. 6), where advection losses become dominant. In the range \( z > 1 \) pc, \( R_{sp} \) is large enough to suppress the effect of photon-photon absorption. In the case of \( L_j = 1.2 \times 10^{46} \text{ erg s}^{-1} \), the synchrotron and IC emission achieve levels of \( \gtrsim 5 \times 10^{39} \text{ erg s}^{-1} \) in hard X-rays and \( \sim 10^{38} \text{ erg s}^{-1} \) in gamma rays, respectively.

In Fig. 4 the bolometric luminosities (synchrotron + IC) at different \( z \) and for a variety of stellar distributions are shown. In Section 6 we have commented about the non-thermal bolometric luminosity \( (L_j + L_{\gamma}) \) produced by the interaction of only one star with the jet at different \( z \), from 1 pc to 1 kpc (maroon-solid lines). However, at \( z \gtrsim r_t \) there are more than one star every time into the jet, and the non-thermal luminosity produced by all the stars into the jet at different \( z \) is also plotted in Fig. 4. In each panel we show the bolometric luminosity produced by the different stellar populations considered in the present study. Note that on the one hand, in the most powerful case (M9-1-0.1) the total bolometric luminosity produced in the jet and in the wind is \( \sim 5 \times 10^{41} \) and \( \sim 10^{39} \text{ erg s}^{-1} \) (y = 2), respectively. On the other hand, in cases with low density of massive stars, the luminosity produced by the cumulative effect of all stars into the jet can be lower than the luminosity produced by only one star interacting with the jet close to z0 if the star formed at \( z \gtrsim r_t \) and migrated close to the jet base.
Figure 9. Spectral energy distribution of the emission up to $z = 1\text{kpc}$ produced by $N_j$ stars inside the jet. The main contributions to the SED are synchrotron radiation and IC scattering; proton-proton interactions are not relevant. Left panel is for the case of $y = 1$, and the right one for $y = 2$. Cases with $M_{bh} = 10^7$ (top), $10^8$ (middle), and $10^9 \text{M}_\odot$ (bottom) are shown.
population of stars, the last two effects soften the emission drop with $z$.

The interaction of only one star with the jet can produce significant amounts of high-energy emission only if the interaction height is below the $z$ at which advection escape dominates the whole particle population. Also, $\sigma_{\text{ap}}$ should be a significant fraction of $\sigma_I$. In this context, we have considered the interaction of a powerful WR star at $z = 1$ pc. The emission produced by IC scattering achieves values as high as $\gtrsim 10^{36}$ erg s$^{-1}$ (considering the contribution of the wind and jet in Fig. 5 in the Fermi range. Such an event would not last long though, about $R_j/\nu_s \sim 300 (R_j/3 \times 10^{17} \text{ cm}) (10^9 \text{ cm s}^{-1}/\nu_s)^{-1} \text{ yr}$. The emission level could be detectable by Fermi only for very nearby sources, like Centaurus A (located at a distance $d \sim 4$ Mpc). The interaction of few WR stars interacting with jets in more distant sources like M87 ($d \sim 16$ Mpc) could be detectable by the forthcoming CTA. The interaction of a star even more powerful than a WR, like a Luminous Blue Variable, may provide $R_{\text{ap}} \sim R_j$, making available the whole jet luminosity budget for particle acceleration.

In the middle/end part of the jet, the interaction of many massive stars can also produce a significant amount of gamma rays. The resulting SED integrated along the whole jet strongly depends on the number of stars inside it. We have considered a Salpeter initial mass function of stars distributed following a power-law spatial distribution (Eq. (4)). In the case of $M_{\text{ini}} = 10^3 M_\odot$, and high accretion rates ($\eta_{\text{accc}} = 1$), gamma-ray luminosities $\sim 10^{38}$ and $5 \times 10^{38}$ erg s$^{-1}$, for $y = 1$ and 2, respectively, may be achieved (see Fig. 5). However, note that few WR inside the jet could actually dominate over the whole main-sequence OB star population.

Although jet/star interactions are very sporadic near the base of the jet, we note that at $z < 1$ pc, clouds from the BLR can also interact with the jet, leading to significant gamma-ray radiation (Araudo et al. 2010). The produced emission in BLR clouds interacting with jets has a stronger dependence on $L_j$ than in the case of stellar winds, because clouds do not have winds and their cross section does not get adjusted to ram pressure balance. Thus, jet/BLR cloud interactions could be more relevant in sources like FR II galaxies.

An interesting (similar) scenario is the interaction of a star forming region (SFR) with the jet. There is evidence that SFRs are located in the torus of some AGN (starburst galaxies), at distances $\sim 100$ pc from the nucleus. In addition, hints of SFRs located in the nuclear region of AGN are also found in galaxies with IR nuclear excess. These galaxies are called nuclear starburst galaxies. The number of OB-type stars in SFRs can be as high as $\sim 10^4$, distributed in a small volume of $\sim 10$ pc$^3$. Then, if one of these compact SFRs interact with the jet at $z \sim 10$ pc, the total luminosity could reach detectable levels, with the resulting radiation presenting rich and complex features. Furthermore, the jet passing through the intra-cloud rich medium can have interesting consequences in the SFR evolution. This scenario will be analyzed in detail in a following paper.

It is noteworthy that, for $\eta_{\text{accc}} \lesssim 1$, one expects $\sim 10^4$ massive stars up to $\sim 1$ kpc. Moreover, as shown in Sect. 3.1, the shocked stellar wind will efficiently mix with the jet. Assuming an average $\dot{M}_w \sim 10^{-6} M_\odot$ yr$^{-1}$, one can estimate the power required to accelerate this mass to the jet Lorentz factor, $\Gamma \dot{M}_w c^2 \sim 6 \times 10^{34}$ erg s$^{-1}$. Despite this is just an order of magnitude estimate, this power tells us that the dynamics of jets with similar or smaller power, i.e. $\lesssim 10^{35}$ erg s$^{-1}$, will be significantly affected by wind mass-loading (e.g. Hubbard & Blackman 2006). Therefore, early-type stars, as low-mass ones (Komissarov 1994), cannot be neglected when studying jet propagation and evolution in galaxies with moderately high star formation. Even for $\eta_{\text{accc}} \sim 0.01$, jets with $L_j \sim 10^{42}$ erg s$^{-1}$ may be strongly affected by the entrainment of wind material (see also the discussion on mass-load in Bosch-Ramon et al. (2012)).

Finally, we remark that since jet-star emission should be rather isotropic (as in all the cases of jet-obstacle interactions), it would be masked by jet beamed emission in blazar sources. Misaligned sources however do not display significant beaming, and for those cases jet-star interactions may be a dominant gamma-ray production mechanism. In the context of AGN unification (e.g. Urry & Padovani 1993), the number of non-blazar AGN should be much larger than that of blazars with the same $L_j$. Close and powerful sources could be detectable by deep enough observations of the Fermi gamma-ray satellite. After few-year exposure, a significant signal from jet-star interactions could be found, and their detection would shed light not only on the jet properties but also on the stellar populations in the vicinity of AGN. The same applies to stars with powerful winds penetrating the jet at its innermost regions, which may be seen as occasional, transient month-scale gamma-ray events.

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