A magnetic reconnection model for explaining the multi-wavelength emission of the microquasars Cyg X-1 and Cyg X-3

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ABSTRACT
Recent studies have revealed that cosmic ray acceleration by a first-order Fermi process in magnetic reconnection current sheets can be rather efficient in the surrounds of compact sources. In this work, we discuss this acceleration mechanism operating in the core region of galactic black hole binaries (or microquasars) and show the conditions under which this can be more efficient than shock acceleration. In addition, we compare the corresponding acceleration rate with the relevant radiative loss rates obtaining the possible energy cut-off of the accelerated particles and also compute the expected spectral energy distribution (SED) for two sources of this class, namely Cygnus X-1 and Cygnus X-3, considering both leptonic and hadronic processes. The derived SEDs are comparable to the observed ones in the low and high energy ranges. Our results suggest that hadronic non-thermal emission due to p-photon interactions may produce the very high energy gamma-rays in these microquasars.

Key words: Microquasars- cosmic ray acceleration: magnetic reconnection- radiation mechanisms: non-thermal.

1 INTRODUCTION
Detected non-thermal radio to gamma-ray emission from galactic binary systems hosting stellar mass black holes, also denominated black hole binaries (BHBs), microquasars, or simply μQSOs (Mirabel & Rodriguez 1994; Tingay et al. 1995; Hjellming & Rupen 1995), provide clear evidence of the production of relativistic particles in their jets and probably also in the innermost regions very close to the black hole (BH). Currently, more than a dozen microquasars have been detected in the galaxy (e.g., Ribo 2002; Paredes & Marti 2003; Zhang 2013).

Usually, the observed radio and infra-red (IR) emission in microquasars is interpreted as due to synchrotron radiation produced by relativistic particles in the jet outflow, while the soft X-ray emission is commonly attributed to the heating of the disk accretion gas as it falls into the BH and the hard X-rays, though still debatable, are related to inverse Compton scattering of seed photons at the base of the compact jet by relativistic electrons (see e.g., Corbel, et al. 2002; Remillard & McClintock 2006; Muno et al. 2001).

Generally, these sources are far from being stable and individual systems have often complex emission structure.

Nevertheless, there are common features to all classes of BHBs and, regarding their X-ray emission (2-100 keV), they show basically two major states: a quiescent and an outburst state (e.g., Remillard & McClintock 2006). The former is characterized by low X-ray luminosities and hard non-thermal spectra. Usually, transient BHBs exhibit this state for long periods, which allows one to obtain the physical parameters of the system. On the other hand, the outburst state corresponds to intense activity and emission, and can be sub-classified in three main active and many intermediary states. According to Remillard & McClintock 2006 (see also Zhang 2013), the three main active states are the thermal state (TS), the hard state (HS) and the steep power law state (SPLS). These states are usually explained as changes in the structure of the accretion flow, as remarked before. During the TS, the soft X-ray thermal emission is believed to come from the inner region of the thin accretion disk that extends until the last stable orbits around the black hole. On the other hand, during the HS the observed weak thermal component suggests that the disk has been truncated at a few hundreds/thousands gravitational radii. The hard X-ray emission measured during this state is dominated by a power-law (PL) component and is often attributed to inverse Compton scattering of soft photons from the outer disk by relativistic electrons in the hot inner region of the
system (e.g., Remillard & McClintock 2006, Malzak et al. 2006). The SFLS is almost a combination of the above two states, but the PL is steeper (Remillard & McClintock 2006, Zhang 2013).

More recently a few of these sources have been also detected in the gamma-ray range with AGILE (Tavani et al. 2009; Bulgarelli et al. 2010; Sabatini et al. 2010a,b, 2013), Fermi-LAT (Atwood et al. 2009; Bodaghee 2013) and MAGIC (Lorentz 2004). For Cyg X-1, for instance, upper limits with 95% confidence level have been obtained in the range of \( \geq 150 \) GeV (Albert et al. 2007), while in the case of Cyg X-3, upper limits of integrated gamma-ray flux above 250 GeV have been inferred by Aleksic et al. (2010). Upper limits in the 0.1-10 GeV range have been also suggested for GRS 1915+105 and GX 339-4.

Several models have been proposed to explain the high energy emission (HE) in microquasars, especially for Cygnus X-1 (Cyg X-1) and Cygnus X-3 (Cyg X-3). For instance, Romero et al. (2003) assumed that the gamma-ray emission is produced in a hadronic jet as a result of the decay of neutral pions created in photon-ion collisions. An alternative model developed by Bosch-Ramon et al. (2005b) assumed that relativistic protons produced in the jet may diffuse through the interstellar medium (ISM) and then interact with molecular clouds and produce gamma-rays out of p – p interactions via neutral pion decay. Another model has been proposed by Piano et al. (2012) in which both, leptonic (via inverse Compton) and hadronic (via neutral pion decay) might account for the observed gamma-ray emission.

All models above postulate that the primary relativistic particles (electrons and protons) are produced by a first-order Fermi mechanism behind shocks in the jet outflow (e.g., Bell 2013; Gallant & Achterberg 1999; Achterberg et al. 2001).

An alternative mechanism has been explored first in the context of microquasars by de Gouveia Dal Pino & Lazarian (2005), hereafter GL05 and later extended also to the framework of AGNs by de Gouveia Dal Pino et al. (2010a, 2010b; de Gouveia Dal Pino & Kowal 2013 and Kadowaki & de Gouveia Dal Pino 2014), in which particles can be accelerated in the surrounds of the BH of these sources, at the jet basis, by a first-order Fermi process occurring in magnetic reconnection sites (current sheets) rather than shocks. These authors found that the energy power extracted from these magnetic reconnection events can be more than sufficient to accelerate the particles and produce the associated radio synchrotron radiation at the nuclear regions of these sources. They also verified that there is a correlation between the magnetic energy power released in these reconnection events and the black hole (BH) mass spanning 10\(^9\) orders of magnitude that can explain the observed radio luminosity of nuclear outbursts from microquasars to low luminous AGNs (de Gouveia Dal Pino et al. 2010a, Kadowaki & de Gouveia Dal Pino 2014), in consistency with the so called fundamental plane (Merloni, Heinz & di Matteo 2003). In addition these authors argued that this process could be related to the X-ray outburst states seen in microquasars as described above.

Also, in recent work the acceleration mechanism above has been successfully tested numerically by means of 3D MHD simulations involving the injection of thousands of test particles in magnetic reconnection domains (Kowal, de Gouveia Dal Pino & Lazarian 2011, 2012). The process allows for an exponential growth with time of the particle kinetic energy to relativistic values and the development of a power law spectrum which is in consistency with the observations.

In this work, we investigate the origin of the non-thermal emission of the microquasars Cyg X-1 and Cyg X-3 assuming the model and the acceleration mechanism above applied to the nuclear region of these sources aiming at reproducing their observed spectral energy distribution (SED) from radio to gamma-rays during outburst states. We consider the relevant radiative loss mechanisms due to the interactions of the accelerated particles with the ambient matter, magnetic and radiation fields and also assess the importance of the acceleration by magnetic reconnection in comparison to shock acceleration.

In Section 2 we describe in detail our acceleration model. In Section 3 we describe the equations employed to calculate the emission processes from radio to gamma-ray energies. In Sections 4 and 5 we show the results of the application of the acceleration and emission model to Cyg X-1 and Cyg X-3, respectively; and in Section 6 we discuss our results and draw our conclusions.

## 2 OUR PARTICLE ACCELERATION SCENARIO

We assume here that relativistic particles are accelerated in the core of the microquasar in the surrounds of the BH, near the basis of the jet launching region, as a result of events of fast magnetic reconnection. This acceleration model has been described in detail in earlier work (de Gouveia Dal Pino & Lazarian 2005; de Gouveia Dal Pino et al. 2010a, 2010b; Kadowaki & de Gouveia Dal Pino 2014) and we summarize here its main assumptions. We consider a magnetized standard (geometrically thin and optically thick) accretion disk around a BH as in the cartoon of Fig. 1. A magnetosphere around the central BH can be established from the drag of magnetic field lines by the accretion disk. The large-scale poloidal magnetic field in the disk corona can in turn be formed by the action of a turbulent dynamo inside the accretion disk (see GL05 and de Gouveia Dal Pino et al. 2010a).
Magnetic reconnection: a CR accelerator

and references therein). This poloidal magnetic flux summed to the disk differential rotation gives rise to a wind that removes angular momentum from the system and increases the accretion rate. This will increase the ram pressure of the accreting material that will then squeeze the magnetic lines in the inner disk region and press them against the lines anchored into the BH horizon allowing for a fast magnetic reconnection event to occur (see the Y shaped zone in Fig. 1). As shown in GL05 (and de Gouveia Dal Pino et al. 2010a), when the accretion rate reaches values close to the Eddington limit the magnetic reconnection event releases substantial magnetic energy. The total magnetic power released by these reconnection events in the torus around the BH as shown in Figure 1, is given by (de Gouveia Dal Pino & Lazarian 2005; de Gouveia Dal Pino et al. 2010a; Kadowaki & de Gouveia Dal Pino 2014):

\[ W = 8.67 \times 10^{33} \alpha^{-0.62} \beta^{0.56} M^{0.59} R_X^{0.78} l_X^{0.69} \text{ erg s}^{-1} \] (1)

where 0.05 ≤ α < 1 is the Shakura-Sunyaev disk viscosity parameter which we here assume to be of the order of 0.5, ξ is the mass accretion disk rate in units of the Eddington rate (\( \xi = \dot{M}/M_{\dot{M}_{\text{Edd}}} \)), which we assume to be ξ ≈ 0.7, M is the BH mass, \( R_X \) is the inner radius of the accretion disk (here taken as \( R_X = 3R_S \), where \( R_S \) is the BH Schwartzchild radius), and \( l_X \) is the length of the reconnection zone (as shown in Figure 1; see also Tables 1 and 2).1

The magnetic power above heats the surrounding gas and accelerates particles. We assume that approximately 50% of the reconnection power is used to accelerate the particles. This is consistent with recent plasma laboratory experiments of particle acceleration in reconnection sheets (e.g., Yamada et al. 2014).2

As in shock acceleration where particles confined between the upstream and downstream flows undergo a first-order Fermi mechanism, GL05 proposed that a similar mechanism would occur when particles are trapped between the two converging magnetic flux tubes moving to each other with \( V_R \) in a magnetic reconnection current sheet. They showed that, as particles bounce back and forth due to head-on collisions with magnetic fluctuations in the current sheet, their energy after a round trip increases by \( < \Delta E/E > \approx 8V_R/c \), which implies a first-order Fermi process with an exponential energy growth after several round trips (see also de Gouveia Dal Pino & Kowal 2013 2014). Under conditions of fast magnetic reconnection which can be induced by anomalous resistivity or by the presence of turbulence in the current sheet (Lazarian & Vishniac 1999), \( V_R \) is of the order of the local Alfvén speed \( V_A \). In particular, at the surroundings of relativistic sources \( V_R \approx V_A \approx c \) and thus the mechanism can be rather efficient (Giannios 2010).

As remarked earlier, this mechanism has been successfully tested by means of numerical simulations in which charged thermal particles were accelerated to relativistic energies into 3D MHD domains of magnetic reconnection (Kowal, de Gouveia Dal Pino & Lazarian 2011, 2012 and de Gouveia Dal Pino & Kowal 2013). Also, tests performed in collisionless fluids by means of PIC simulations have achieved similar results (e.g., Zenitani & Hoshino 2001 Zenitani, Hesse & Klimas 2009; Drake et al. 2006, 2010; Cerutti et al. 2013 2014).

The numerical simulations by Kowal, de Gouveia Dal Pino & Lazarian (2012) indicate that under fast magnetic reconnection the particle acceleration rate within the current sheet is (see also de Gouveia Dal Pino & Kowal 2013; del Valle et al. 2014):

\[ \dot{E}_{\text{acc},M,R,,p}^{-1} = \frac{1}{\alpha_0} \left( \frac{E}{E_0} \right)^{-\alpha_0} t_0^{-1}, \] (2)

where 0.36 ≤ \( \alpha_0 \) ≤ 0.4, \( E_0 = m_p c^2 \), \( t_0 = L_{\text{acc}}/V_A \) which is the Alfvén time scale, with \( V_A = (B/4\pi \rho)^{1/2} \) where B is the local magnetic field and \( \rho \) the fluid density, and \( L_{\text{acc}} \) is the length scale of the acceleration region. Although this result was found from numerical simulations employing protons as test particles, we employ a similar rate for the electrons:

\[ \dot{E}_{\text{acc},M,R,,e}^{-1} = \frac{1}{\alpha_0} \left( \frac{E}{E_0} \right)^{-\alpha_0} t_0^{-1} \sqrt{\frac{m_p}{m_e}}. \] (3)

where \( m_p \) and \( m_e \) are the rest mass of proton and electron respectively. The equations above will be used to compute the acceleration rate in our model in the next sections.

The accelerated particles develop a power law energy distribution (see also Appendix A):

\[ Q(E) \propto E^{-p}, \] (4)

we assume for the power law index a value \( p = 1.8 \). This hard power law index is compatible with analytical predictions (e.g., Drury 2012) and also with the values derived from numerical simulations of particle acceleration by first-order Fermi process in MHD domains of magnetic reconnection (see de Gouveia Dal Pino & Kowal 2013 2014; del Valle et al. 2014) and also PIC simulations in collisionless reconnection domains (Zenitani & Hoshino 2001 Zenitani, Hesse & Klimas 2009; Cerutti et al. 2013 2014).

As stressed in GL05, a diffusive shock may also develop in the surrounds of the magnetic reconnection zone, at the jet launching region, due to the interaction of plasmons and the heated plasma by the magnetic energy released in the fast reconnection event. A similar picture has been also suggested by e.g., Romero, Vieyro & Villa 2010. In this case,
the shock acceleration rate for a particle of energy $E$ in a magnetic field $B$, is given by, (e.g., Spruit 1988):
\[
t_{\text{acc, shock}}^{-1} = \frac{\eta c B}{E},
\]
where $0 < \eta \ll 1$ is a parameter that characterizes the efficiency of the acceleration. We fix $\eta = 10^{-2}$, which describes the efficient acceleration by shocks with velocity $v_s \approx 0.1c$ in the Bohm regime (Romero, Vieyro & Villa 2010).

The accelerated particles lose their energy radiatively via interactions with the surrounding magnetic field (producing synchrotron emission), the photon field (producing inverse Compton, synchrotron-self-Compton, and proton-photon interactions) and with the surrounding matter (producing proton-proton interactions and relativistic Bremsstrahlung radiation).

The ambient magnetic field in the surrounds of the BH calculated from GL05’s model is given by (see also de Gouveia Dal Pino et al. 2010a; Kadowaki & de Gouveia Dal Pino 2014):
\[
B \approx 3.4 \times 10^4 \xi^{1/2} \dot{M}^{0.25} \alpha^{-0.75} \text{ G}
\]
The particle density in the coronal region in the surrounds of the BH is
\[
n_c \approx 8.13 \times 10^{20} \alpha^{-0.5} \xi^{0.75} \dot{M}^{0.12} R_X^{-0.37} l_X^{-0.75} \text{ cm}^{-3}.
\]
The equations above will be employed in Sections 4 and 5 to model the acceleration in the core region of the microquasars Cyg X-1 and Cyg X-3. The volume of the acceleration region in our model is taken to be the cylindrical shell where magnetic reconnection takes place, as in Figure 1. This shell has a length $l_X$, and inner and outer radii $R_X$ and $R_X + \Delta R_X$, respectively, so that the volume of the acceleration zone is $2\pi l_X (R_X + \Delta R_X)$, where $\Delta R_X$ is the width of the current sheet given by (GL05 and de Gouveia Dal Pino et al. 2010a):
\[
\Delta R_X \approx 1.5 \times 10^{-5} \alpha^{-0.62} \xi^{0.43} \dot{M}^{0.1} R_X^{0.28} l_X^{0.68} \text{ cm}.
\]
In sections 4 and 5, we will also need the accretion disk temperature in order to evaluate its black body radiation field:
\[
T_d \approx 2.67 \times 10^5 \alpha^{-0.25} \dot{M}^{0.12} R_X^{0.37} \text{ K}.
\]

In the following section we discuss the relevant radiative loss processes for electrons and protons which will allow the construction of the SED of these sources for comparison with the observations.

3 EMISSION AND ABSORPTION MECHANISMS

3.1 Interactions with matter

Charged particles with energy $E$, mass $m$ and charge number $Z$ spiralling in a magnetic field $\vec{B}$ emit synchrotron radiation at a rate
\[
t_{\text{synch}}^{-1}(E) = \frac{4}{3} \left( \frac{m_e}{m} \right)^3 \frac{\sigma_T B^2}{m_e c^5 \alpha_f c} \frac{E}{mc^2},
\]
where $m_e$ is the electron mass and $\sigma_T$ is the Thompson cross section. The synchrotron spectrum radiated by a distribution of particles $N(E)$ (see appendix A) as function of the scattered photon energy ($E_\gamma$) (in units of power per unit area) is
\[
L_{\gamma}(E_\gamma) = \frac{E_\gamma V}{4\pi d^2} \int_{E_{\text{min}}}^{E_{\text{max}}} dE N(E) \frac{E_\gamma}{E} \int_{\frac{E}{E_\gamma}}^{\infty} K_{5/3}(\xi) d\xi,
\]
where $V$ is the volume of the emission region, $d$ is the distance of the source from us, $h$ is the Planck constant, $K_{5/3}(\xi)$ is the modified Bessel function of 5/3 order, and the characteristic energy $E_c$ is
\[
E_c = \frac{3 c^3 B}{8 \pi m_e c^2} \left( \frac{E}{m_e c^2} \right)^2.
\]
In these calculations we assumed that the particle velocity is perpendicular to the magnetic field.

To compute equation (11) we used the approximation
\[
x \int_{2}^{\infty} K_{5/3}(\xi) d\xi \approx 1.85 x^{1/3} e^{-x}.
\]
Practically, the synchrotron emission of the electrons dominates the low energy photon background which is a proper target for both inverse Compton (IC) and $pr\gamma$ interactions (see below; see also Reynoso, Medina & Romero 2011).

The number density of multi-wavelength synchrotron scattered photons (in units of energy per volume), has been approximated as (Reynoso, Medina & Romero 2011; Zhang, Chen & Fang 2008; Romero, del Valle & Orellana 2010)
\[
n_{\text{synch}}(\epsilon) = \frac{L_{\gamma}(\epsilon)}{c E_v} \frac{\epsilon}{4\pi d^2},
\]
where $r_X$ stands for the radius of the emission region and $\epsilon$ for the scattered synchrotron radiation energy. More precisely, $\epsilon$ corresponds to the photon energy of the multi-wavelength target radiation field for synchrotron self Compton (SSC) and $pr\gamma$ interactions. The volume $V$ of the emission region in our model is taken as the spherical region that encompasses the cylindrical shell where magnetic reconnection particle acceleration takes place in Figure 1. Considering that the torus extends up to $l_X$, then $r \approx l_X$ and the effective emission zone in our model has an approximate volume $4\pi l_X^3/3$ (see Tables 1 and 2).

3.2 Interactions with matter

3.2.1 Bremsstrahlung

When a relativistic electron accelerates in the presence of the electrostatic field of a charged particle or a nucleus of charge Ze then Bremsstrahlung radiation is produced. For a completely ionized plasma with ion number density $n_i$, the Bremsstrahlung cooling rate is (Berezhinskii 1990; Romero, Vieyro & Villa 2010)
\[
t_{\text{Br}}^{-1} = 4n_i Z^2 \alpha_f c \left[ \ln \left( \frac{2E_e}{m_e c^2} \right) - \frac{1}{3} \right],
\]
where $r_0$ is the electron classical radius and $\alpha_f$ stands for the fine structure constant. The relativistic Bremsstrahlung luminosity (in units of power per unit area) is given by (Romero, del Valle & Orellana 2010)
\[
L_{\gamma}(E_\gamma) = \frac{E_\gamma V}{4\pi d^2} \int_{E_{\gamma}}^{\infty} n_e c E_\gamma dE_\gamma,
\]
where

\[ \sigma_B(E_\gamma, E_\gamma) = \frac{4a_f r_p^2 \Phi(E_\gamma, E_\gamma)}{E_\gamma}, \]

and

\[ \Phi(E_\gamma, E_\gamma) = \left[ 1 + \left( 1 - \frac{E_\gamma}{E_H} \right)^2 - \frac{2}{3} \left( 1 - \frac{E_\gamma}{E_H} \right) \right] \times \left[ \ln \frac{2E_\gamma(E_\gamma - E_{\pi})}{mc^2E_\gamma} - \frac{1}{2} \right]. \]

### 3.2.2 pp interactions

One relevant gamma-ray production mechanism is the decay of neutral pions which can be created through inelastic collisions of the relativistic protons with nuclei of the corona that surrounds the accretion disk. In this case the cooling rate is given by (Kelner 2009)

\[ t_p^{-1} = n_i c \sigma_{pp} k_{pp}, \]

where \( k_{pp} \) is the total inelasticity of the process of value \( \sim 0.5 \). The corresponding cross section for inelastic pp interactions \( \sigma_{pp} \) can be approximated by

\[ \sigma_{pp}(E_p) = \left( 34.3 + 1.88L + 0.25L^2 \right)
\times \left[ 1 - \left( \frac{E_{th}}{E_p} \right)^4 \right]^2 \text{mb,} \]

where \( L = \ln \left( \frac{E_p}{mc^2} \right) \), and the proton threshold kinetic energy for neutral pion \((\pi^0)\) production is \( E_{th} = 2m_e c^2(1 + \frac{m_e^2}{E_p^2}) \approx 280\text{ MeV} \), while \( m_e c^2 = 134.97\text{ MeV}\) is the rest energy of the \( \pi^0 \) (Villa & Aharonian 2009). This particle decays in two photons with a probability of 98.8%.

The luminosity can be calculated by

\[ L_\gamma(E_\gamma) = \frac{E_\gamma^2 V}{4\pi d^2} \Gamma(E_\gamma), \]

where \( \Gamma(E_\gamma) = (E_\gamma c^{-1}\text{cm}^{-3}\text{s}^{-1}) \) is the gamma-ray emissivity.

For proton energies less than 0.1 TeV, \( \Gamma(E_\gamma) \) is

\[ \Gamma(E_\gamma) = 2 \int_{E_{min}}^{\infty} \frac{q_\pi(E_\pi)}{\sqrt{E_\pi^2 - m_e^2c^4}} dE_\pi, \]

where \( E_{min} = E_\gamma + m_e^2 c^4/4E_\gamma \) and \( q_\pi(E_\pi) \) is the pion emissivity. An approximate expression for \( q_\pi(E_\pi) \) can be calculated using the \( \delta \)-function (Aharonian & Atoyan 2000).

For this purpose, the neutral pion takes a fraction \( k_\pi \) of the kinetic energy of the proton \( E_{kin} = E_p - m_p c^2 \) (Villa & Aharonian 2009). The pion emissivity is then given by

\[ q_\pi(E_\pi) = \frac{n_i}{k_\pi} \int \delta(E_\pi - k_\pi E_{kin}) \sigma_{pp}(E_p) N_p(E_p) dE_p \]

\[ = \frac{n_i}{k_\pi} \sigma_{pp}(m_p c^2 + E_\pi/k_\pi) N_p(m_p c^2 + E_\pi/k_\pi). \]

The target ambient nuclei density is given by \( n_i \) and \( N_p(E_p) \) stands for the energy distribution of the relativistic protons.

For proton energies in the range GeV-TeV, \( k_\pi \approx 0.17 \) (Gaisser 1990), the total cross section \( \sigma_{pp}(E_p) \) can be approximated by

\[ \sigma_{pp}(E_p) \approx \begin{cases} 30 \left[ 0.95 + 0.06 \ln \left( \frac{E_{kin}}{1\text{ TeV}} \right) \right] \text{mb} & \text{if } E_{kin} \geq 1\text{GeV}, \\ 0 & \text{if } E_{kin} < 1\text{GeV}. \end{cases} \]

For proton energies greater than 0.1 TeV, the gamma-ray emissivity is

\[ q_\pi(E_\pi) = \frac{c n_i}{E_p} \int_{E_{kin}}^{\infty} \sigma_{\pi\pi}(E_p) N_p(E_p) \frac{dE_p}{E_p} \]

\[ = \frac{c n_i}{E_p} \int_{0}^{1} \sigma_{\pi\pi}(E_\pi x) N_p(E_\pi x) F(x) \frac{dx}{x}. \]

The inelastic pp cross section is approximately given by

\[ \sigma_{\pi\pi}(E_p) = \left( 34.3 + 1.88 L + 0.25 L^2 \right)[1 - \left( \frac{E_{th}}{E_p} \right)^4]^2 \text{mb}, \]

Here \( E_{th} = m_p + 2m_e = \frac{m^2}{2m_p} = 1.22 \text{ GeV} \) is the threshold energy of the proton to produce neutral pions \( \pi^0 \) and the number of photons with energy in the interval \((x, x + dx)\) (where \( x = E_\gamma/E_p \)) which is created per pp collision can be approximated by (Villa & Aharonian 2009)

\[ F_\gamma(x, E_p) = B_\gamma \left( \frac{x}{x_0} \right)^{1 - x_0^4} \]

\[ \times \left[ \frac{1}{(1 - x_0^4)} \frac{4\beta x_0^4}{(1 + x_0^4) - 4k_\gamma x_0^4(1 - 2x_0^4)} \right]. \]

The best least-squares fit to the numerical calculations yield:

\[ B_\gamma = 1.30 + 0.14L + 0.011L^2, \]

\[ \beta_0 = (1.79 + 0.11L + 0.008L^2)^{-1}, \]

\[ k_\gamma = (0.801 + 0.049L + 0.014L^2)^{-1}. \]

Where \( L = \ln(E_p/1\text{TeV}) \) and \( 0.001 \leq x \leq 0.1 \) for details see Villa & Aharonian 2009.

### 3.3 Interactions with the radiation field

Energetic electrons transfer their energy to low energy photons causing them to radiate at high energies (inverse Compton process). On the other hand, when high energy protons interact with low energy photons (\( p - \gamma \) interactions) they produce pions and gamma-ray photons with energies larger than \( 10^8 \text{ eV} \) in the so called photomeson process.

#### 3.3.1 Inverse Compton

The IC cooling rate for an electron in both Thomson and Klein-Nishina regimes is given by (Blumenthal & Gould 1970)

\[ \dot{E}_e^{-1} = \frac{1}{E_e} \int_{E_{min}}^{E_{max}} \int_{E_{ph}}^{E_{th}} \left( \frac{E_\gamma}{E_p} - E_{ph} \right) \frac{dE_\gamma}{ddE_\gamma} dE_\gamma. \]

Here \( E_{ph} \) and \( E_\gamma \) are the incident and scattered photon energies, and

\[ \frac{dN}{ddE_\gamma} = \frac{2\pi^2 m_p c^2 n_{ph}(E_{ph}) dE_{ph}}{E_p^2} \mathcal{F}(q), \]
where \( n_{ph}(E_{ph}) \) is the target photon density (in the units of energy\(^{-1}\)volume\(^{-1}\)) and

\[
F(q) = 2q\ln q + (1 + 2q)(1 - q)\frac{(\Gamma q)^2}{1 + \Gamma},
\]

\[
\Gamma = 4E_{ph}E_e/(mc^2)^2,
\]

\[
q = \frac{E_{\gamma}}{\Gamma(E_c - E_{\gamma})}.
\]

We consider two target photon fields to model the IC emission. We consider interactions of the accelerated electrons with photons produced in the synchrotron emission, in which case the process is synchrotron-self-compton (SSC, eq. 14), and photons emitted by the surface of the accretion disk, although this second contribution is less important, which we will see later. This photon field can be represented by a black body radiation with the disk temperature and is given by

\[
n_{bb}(E_{ph}) = \frac{1}{\pi^2\lambda_0^2m_e^2c^2}\frac{E_{ph}^2}{(E_{ph}c^2)^2}\left[\frac{1}{\exp(\frac{E_{ph}}{\lambda_0c})} - 1\right].
\]

Here \( \lambda_0, t \) and \( k \) are Compton wavelength, disk temperature and Boltzmann constant respectively.

Taking into account the Klein-Nishina effect on the process of a proton with energy \( E_{\gamma} \), the spectrum of photons produced can be calculated from

\[
\frac{dL_{IC}(E_{\gamma})}{dE_{\gamma}} = \frac{E_{\gamma}^2}{4\pi\Gamma^2} \int_{E_{\gamma}}^{E_{\gamma,max}} dE_e N_e(E_e)
\]

\[
\times \int_{E_{ph,min}}^{E_{ph,max}} dE_{ph} P_{IC}(E_{\gamma}, E_{ph}, E_e),
\]

where \( P_{IC}(E_{\gamma}, E_{ph}, E_e) \) is the spectrum of photons scattered by an electron of energy \( E_e = \gamma_e mc^2 \) in a target radiation field of density \( n_{ph}(E_{ph}) \), according to Blumenhaw & Gould (1970), it is given by

\[
P_{IC}(E_{\gamma}, E_{ph}, E_e) = \frac{3\sigma_C(m_e c^2)^2}{4E_{\gamma}^2} \frac{n_{ph}(E_{ph})}{E_{ph}} F(q),
\]

and for the scattered photons there is a range which is

\[
E_{ph} \leq E_{\gamma} \leq \frac{\Gamma}{1 + \Gamma} E_e.
\]

3.3.2 Photomeson production \((p - \gamma)\)

The photomeson production takes place for photon energies greater than \( E_{th} \approx 145\text{MeV} \). A single pion can be produced in an interaction near the threshold and then decay giving rise to gamma-rays. In our model the appropriate photons come from the synchrotron radiation.

\(^3\) We note that the contribution of target photons due to the radiation field produced by the companion star is found to be irrelevant in our model (e.g., Bosch-Ramon et al. 2005a).

\(^4\) We find that for photomeson production, the radiation from the accretion disk and from the companion star are irrelevant compared to the contribution from the synchrotron emission.

for this mechanism in an isotropic photon field with density \( n_{ph}(E_{ph}) \) can be calculated by Stecker (1965):

\[
t_{\gamma\gamma}^{-1}(E_{\gamma}) = \frac{e}{2\gamma_{\gamma}^2} \int_{E_{th}}^{E_{\gamma}} dE_{\gamma} \frac{n_{ph}(E_{ph})}{E_{ph}}
\]

\[
\times \int_{E_{th}}^{E_{\gamma}} dE_{\gamma} \sigma_{\gamma\gamma}(E_{\gamma}) K_{p\gamma}(E_{\gamma}) \epsilon_{\gamma},
\]

where \( \gamma_{\gamma} = \frac{E_{\gamma}}{m_e c^2}, \epsilon_{\gamma} = \text{the photon energy in the rest frame of the proton and } K_{p\gamma} \) is the inelasticity of the interaction. Atoyan & Dermer (2003) proposed a simplified approach to calculate the cross-section and the inelasticity which are given by

\[
\sigma_{p\gamma}(\epsilon_{\gamma}) \approx \begin{cases} 340 \text{ mbarn} & 300 \text{ MeV} \leq \epsilon_{\gamma} \leq 500 \text{ MeV}, \\ 120 \text{ mbarn} & \epsilon_{\gamma} > 500 \text{ MeV}, \end{cases}
\]

and

\[
K_{p\gamma}(\epsilon_{\gamma}) \approx \begin{cases} 0.2 & 300 \text{ MeV} \leq \epsilon_{\gamma} \leq 500 \text{ MeV}, \\ 0.6 & \epsilon_{\gamma} > 500 \text{ MeV}. \end{cases}
\]

To find the luminosity from the decay of pions, we use the analytical approach proposed by Atoyan & Dermer (2003). Taking into account that each pion decays into two photons, the \( p - \gamma \) luminosity is

\[
L_{p\gamma}(E_{\gamma}) = \frac{2E_{\gamma}^2}{4\pi d^2} \int \frac{Q_{\gamma\gamma}^{(p\gamma)}}{E_{\gamma}} \delta(E_{\gamma} - 0.5E_{\gamma}) dE_{\gamma}
\]

\[
= 20 \frac{E_{\gamma}^2}{4\pi d^2} N_P(10E_{\gamma}) \omega_{p\gamma,\gamma}(10E_{\gamma}) n_{\gamma0}(10E_{\gamma}),
\]

where \( Q_{\gamma\gamma}^{(p\gamma)} \) is the emissivity of the neutral pions given by

\[
Q_{\gamma\gamma}^{(p\gamma)} = 5N_P(5E_{\gamma}) \omega_{p\gamma,\gamma}(5E_{\gamma}) n_{\gamma0}(5E_{\gamma}),
\]

\( \omega_{p\gamma} \) stands for the collision rate which is

\[
\omega_{p\gamma}(E_{p}) = \frac{m_p^2 c^2}{2E_p^2} \int_{E_{th}}^{\infty} dE_{\gamma} \frac{n_{ph}(E_{ph})}{E_{ph}} \int_{E_{th}}^{2E_{ph}\gamma_{\gamma}} dE_{\gamma} \sigma_{p\gamma}(E_{\gamma}) \epsilon_{\gamma},
\]

and \( n_{\gamma0} \) is the mean number of neutral pions produced per collision given by

\[
n_{\gamma0}(E_p) = 1 - P(E_p) \xi_{pn}.
\]

In the single-pion production channel, the probability for the conversion of a proton to a neutron with the emission of a \( \pi^+ - meson \) is given by \( \xi_{pn} \approx 0.5 \). For photomeson interactions of a proton with energy \( E_p \), the interaction probability is represented by \( P(E_p) \), which is

\[
P(E_p) = \frac{K_2 - K_{p\gamma}(E_p)}{K_2 - K_1}
\]

\[
K_{p\gamma} = \frac{1}{t_{p\gamma}(\gamma_{\gamma}) \omega_{p\gamma}(E_p)}
\]

The inelasticity in the single-pion channel is approximated as \( K_1 \approx 0.2 \), whereas \( K_2 \approx 0.6 \). For energies above 500 MeV the mean inelasticity \( K_{p\gamma} \) is
3.4 Absorption

Gamma-rays can be annihilated by the surrounding radiation field via electron-positron pair creation: \( \gamma + \gamma \rightarrow e^+ + e^- \). In microquasars, besides the radiation field of the tight companion star, coronal and accretion disk photons can also absorb \( \gamma \)-rays. However, it has been shown by [Cerutti et al. 2011] that the absorption due to coronal photons is too small as compared with the contribution from the disk.

Adopting the same absorption model for the disk radiation field of these authors we find the the disk contribution to gamma-ray absorption is less relevant than that of the stellar companion, generally a Wolf-Rayet star, which produces UV radiation. To evaluate the optical depth due to this component, we have adopted the model described by Sierpowska-Bartosik & Torres (2008), (see also Dubus 2006, Zdziarski & Mikolajewska 2013). This process is possible only above a kinematic energy threshold given by

\[ E_\gamma \epsilon (1 - \cos \theta) \geq 2m_e^2c^4, \quad (49) \]

and

\[ E_\gamma \epsilon > (m_e c^2)^2, \quad (50) \]

in head-on collisions (Romero, Vieyro & Villa 2010), where \( E_\gamma \) and \( \epsilon \) are the energies of the emitted gamma-ray and the ambient photons and \( \theta \) is the collision angle in the laboratory reference frame.

The attenuated luminosity \( L_\gamma (E_\gamma) \) after the \( \gamma \)-ray travels a distance \( l \) is (Romero & Christiansen 2005)

\[ L_\gamma (E_\gamma) = L_\gamma^0 (E_\gamma) e^{-\tau (l, E_\gamma)}, \quad (51) \]

where \( L_\gamma^0 \) is the intrinsic coronal gamma-ray luminosity and \( \tau (l, E_\gamma) \) is the optical depth. The differential optical depth is given by:

\[ d\tau = (1 - \mu) n_{\gamma \gamma} \sigma_{\gamma \gamma} d\Omega dl' \quad (52) \]

where \( d\Omega \) is the solid angle of the target soft photons, \( \mu \) is the cosine of the angle between the gamma-ray and the arriving soft photons, \( l' \) is the path along the gamma-ray emission and \( n_{\gamma \gamma} \) is the black-body photon density in \( \text{cm}^{-3} \text{erg}^{-1}\text{sr}^{-1} \).

The \( \gamma \gamma \) interaction cross-section \( \sigma_{\gamma \gamma} \) is defined as (Gould & Schder 1967)

\[ \sigma_{\gamma \gamma} (\epsilon, E_\gamma) = \frac{\pi \beta^2}{2} \left( 1 - \beta^2 \right) \left[ 2\beta (\beta^2 - 2) + (3 - \beta^3) \ln \left( \frac{1 + \beta}{1 - \beta} \right) \right], \quad (53) \]

where \( r_0 \) is the classical radius of the electron and

\[ \beta = 1 - \left( \frac{m_e c^2}{\epsilon E_\gamma} \right)^{1/2}. \quad (54) \]

The companion star with radius \( R_* \) and a black-body surface temperature \( T_* \) produces a photon density at a distance \( d \) from the star

\[ n_{\gamma \gamma} = \frac{2 \epsilon^2}{h^3 c^3 \exp(\epsilon/k T_*)} \frac{R_*^2}{d^4}. \quad (55) \]

In the absorption models proposed by Sierpowska-Bartosik & Torres (2008) and Dubus (2006), the geometrical parameters \( d, \epsilon \), and \( l \) are strongly dependent on the viewing angle \( \theta \) and the orbital phase \( \phi_0 \). In the superior conjunction, the compact object is behind the star and the orbital phase is \( \phi_0 = 0 \). We here consider the same orbital phase that has been observed during the high energy observations for Cyg X-1 and Cyg X-3. (For more details on the geometrical conditions of the binary system and the integration extremes, see Sierpowska-Bartosik & Torres (2008) and Dubus 2006.

We note that the pairs produced by the absorbed gamma-rays may emit predominantly synchrotron emission in the surrounding magnetic fields (Bosch-Ramon et al. 2008), but their emission is expected to be negligible compared to the other synchrotron processes of the system. We thus neglect this effect in our treatment of pair absorption (Zdziarski et al. 2014).

4 APPLICATION TO CYGNUS X-1

Cyg X-1 is a widely studied black hole (BH) binary system [Malyshev et al. 2013] at a distance of 1.86-2.2 kpc [Reid et al. 2011, Zdziarski 2005] which is accreting from a high mass companion star orbiting around the BH with a period of 5.6 days (Gies et al. 2008). The orbit inclination is between 25\(^{\circ}\) and 35\(^{\circ}\) (Gies & Bolton 1986) with an eccentricity of \( \sim 0.018 \) (Orosz et al. 2011), so that one can assume an approximate circular orbit with a radius \( a_{orb} \).

The parameters of the model for Cyg X-1 are tabulated in Table 1. The values for the first eight parameters in the Table have been calculated from eqs. 1 & 6-9 above. We take for the accretion disk inner radius the value \( r_X = 3R_S \), where \( R_S \) is the BH Schwarzschild radius, and for the extension \( bX \) of the reconnection region (see Fig. 1), we consider the value \( bX \simeq 10R_X \) (GL05, de Gouveia Dal Pino et al. 2010a). As remarked in Section 3, the volume \( V \) of the emission region in Table 1 was calculated by considering the spherical region that encompasses the reconnection region in Figure 1. The black hole mass has been taken from Orosz (2011).

Figures 2 and 3 show the cooling rates for the different energy loss processes described in Section 3 (eqs. 10, 15, 19, 31 and 40) for electrons and protons, respectively. These are compared with the acceleration rates due to first-order Fermi acceleration by magnetic reconnection (eqs. 2 & 3) and to shock acceleration (eq. 5). We notice that for both protons and electrons the acceleration is dominated by the first-order Fermi magnetic reconnection process in the core region. Besides, the main radiative cooling process

| \( B \) | Magnetic field (G) | \( 5.4 \times 10^7 \) |
|---|---|---|
| \( n_\gamma \) | Coronal particle number density \( (cm^{-3}) \) | \( 1.5 \times 10^{16} \) |
| \( T_d \) | Disk temperature (K) | \( 1.4 \times 10^7 \) |
| \( W \) | Reconnection power (erg/s) | \( 3.3 \times 10^{44} \) |
| \( R_s \) | Inner radius of disk (cm) | \( 1.3 \times 10^7 \) |
| \( l_X \) | Height of reconnection region | \( 1.3 \times 10^8 \) |
| \( V_{acc} \) | Volume of acceleration region \( (cm^3) \) | \( 1.2 \times 10^{22} \) |
| \( V \) | Volume of emission region \( (cm^3) \) | \( 9.4 \times 10^{34} \) |
| \( d \) | Distance (kpc) | 2 |
| \( M \) | Mass of BH \((M_\odot)\) | 14.8 |
| \( p \) | Particle power index | 1.8 |
| \( R_s \) | Stellar radius (cm) | \( 1.5 \times 10^{12} \) |
| \( T_s \) | Stellar temperature (K) | \( 3 \times 10^4 \) |
| \( r_{orb} \) | Orbital radius (cm) | \( 3.4 \times 10^{12} \) |
| \( \theta \) | Viewing angle | \( \pi/6 \) |
the surrounding radiation field. As stressed, our calculations indicate that this process is dominated by the radiation field of the companion star. As a result, the opacity depends on the phase of the orbital motion and on the viewing angle.

The parameters employed in the evaluation of this absorption are in the last four lines of Table 1, and have been taken from Romero, del Valle & Orellana (2010). It has been proposed from MAGIC observations (Albert et al. 2007) that the gamma-ray production and absorption are maximized near the superior conjunction (Bodaghee 2013) at phase $\phi_b = 0.91$. In our calculations we considered this orbital phase for Cyg X-1.

The calculated opacity according to the equations above results in a very high energy gamma ray absorption. We find that the produced gamma-rays are fully absorbed in the energy range of 30 GeV-0.3 TeV which causes the energy gap seen in the calculated SED in Figure 4. The observed upper limits by MAGIC plotted in the diagram in this range are possibly originated outside the core, along the jet where $\gamma$-ray absorption by the stellar radiation is not important (see also Romero, del Valle & Orellana 2010).

We note that in Figure 4 the observed flux in radio (10 $\mu$eV – 0.1 eV) and soft gamma-ray ($10^5$ – $10^8$ eV) are explained by leptonic synchrotron and IC processes according to the present model. In the range 10 MeV- 0.6 GeV, inverse Compton is the main mechanism to produce the observed data as a result of interactions between the high energy electrons with both synchrotron photons (SSC) and photons emitted from the disk. At energies in the range 0.6 GeV - 3 TeV, neutral pion decays reproduce the observed gamma-rays. These neutral pions result from $pp$ and $p – photon$ interactions. In the range of 1 GeV- 30 GeV,
p-p collisions are the dominant radiation mechanism, but in the very high energy gamma-rays, interactions of relativistic hadrons (mostly protons) with scattered photons from synchrotron radiation may produce the observed flux.

The observed emission in the near infrared (0.1 eV-10 eV), represented in Figure 4 by redish stars, is attributed to thermal blackbody radiation from the stellar companion and the X-ray emission (1 keV-0.1 MeV), also represented in Figure 4 by redish stars, is believed to be due to thermal Comptonization of the disk emission by the surrounding coronal plasma of temperature \( \sim 10^7 \) K (Di Salvo et al. 2001; Zdziarski et al. 2012). For this reason, these observed data are not fitted by the coronal non-thermal emission model investigated here.

## 5 APPLICATION TO CYGNUS X-3

Cyg X-3 is also a high mass X-ray binary that possibly hosts a BH (Zdziarski & Mikolajewska 2013) and a Wolf-Rayet as a companion star (van Kerkwijk et al. 1992). The system is located at a distance of 7-10 kpc (Bonnet-Bidaud & Chardin 1988) and has an orbital period of 4.8 h and an orbital radius \( \approx 3 \times 10^{13} \) cm (Piano et al. 2012).

Our model parameters for Cyg X-3 are given in Table 2. As in Cyg X-1, the values for the first eight parameters were calculated from eqs. 1 & 6-9 which describe the magnetic reconnection acceleration model in the core region. We have also used for the accretion disk inner radius the value \( R_X = 3R_S \) and for the extension \( l_X \) of the reconnection region the value \( l_X = 10R_X \) (GL05; de Gouveia Dal Pino et al. 2010a). The BH mass has been taken from Schmutz, Geballe & Schild (1996). The cooling and acceleration rates for electrons and protons are depicted in Figures 5 and 6, respectively. The maximum electron and proton energies in both diagrams are obtained from the intercept between the acceleration rate curve and the dominant radiative loss rate curve. As in Cyg X-1, it is clear from these that acceleration by magnetic reconnection is dominating over shock acceleration in the core region. Synchrotron emission is the main mechanism to cool the electrons which may reach energies as high as \( \sim 2 \) GeV, while the most important loss mechanism for protons is \( p-\gamma \) interactions with synchrotron photons. They can be accelerated up to \( \sim 0.3 \) PeV.

In this system, the close proximity \( (R_d \approx 3 \times 10^{11} \text{cm}) \), the high stellar surface temperature \( (T_\star \sim 10^5 \text{K}) \), and the high stellar luminosity \( (L_\star \sim 10^{38} \text{erg s}^{-1}) \) of the Wolf-Rayet star may result a considerable attenuation of the gamma-rays via \( \gamma-\gamma \) pair production (Bednarz 2010). The detection of TeV gamma-rays in Cyg X-3, therefore, relies on the competition between the production and the attenuation process above.

Figure 7 shows the calculated SED compared to the observed data for this source. The gamma-ray absorption was calculated from Eq. 52, employing the UV field of the companion star which is a more significant target than the radiation fields of the accretion disk or the corona (see the stellar parameters in the last five lines of Table 2 which were taken from Cherepashchuk & Moffat 1994). The orbital phase considered was \( \phi_b = 0.9 \), near the superior conjunction (Aleksic et al. 2010), as in Cyg X-1. The energy gap

### Table 2. Model parameters for Cyg X-3.

| Parameter | Value |
|-----------|-------|
| \( B \) | Magnetic field (G) | \( 4 \times 10^7 \) |
| \( n_c \) | Coronal particle number density \( (\text{cm}^{-3}) \) | \( 9.8 \times 10^{15} \) |
| \( T_d \) | Disk temperature (K) | \( 1.2 \times 10^7 \) |
| \( W \) | Reconnection power \( (\text{erg/s}) \) | \( 3.6 \times 10^{34} \) |
| \( R_s \) | Inner radius of disk (cm) | \( 2 \times 10^7 \) |
| \( l_X \) | Height of reconnection region | \( 2 \times 10^8 \) |
| \( V_{acc} \) | Volume of acceleration region \( (\text{cm}^3) \) | \( 1.2 \times 10^{22} \) |
| \( V \) | Volume of emission region \( (\text{cm}^3) \) | \( 9.4 \times 10^{24} \) |
| \( d \) | Distance (kpc) | 8 |
| \( M \) | Mass of BH \( (M_\odot) \) | 17 |
| \( p \) | Particle power index | 1.8 |
| \( R_\star \) | Stellar radius (cm) | \( 2 \times 10^{11} \) |
| \( T_\star \) | Stellar temperature (K) | \( 9 \times 10^4 \) |
| \( r_{orb} \) | Orbital radius (cm) | \( 4.5 \times 10^{11} \) |
| \( \theta \) | Viewing angle | \( \pi/6 \) |

Figure 5. Acceleration and cooling rates for electrons in the core region of Cyg X-3.

Figure 6. Acceleration and cooling rates for protons in the core region of Cyg X-3.
caused by this gamma-ray absorption is shown in the figure 7 in the 50GeV − 0.4TeV.

The contributions of \( pp \) and proton-photon interactions are the dominant ones in the high energy gamma-ray range. These processes become more relevant in the coronal region around the BH since the magnetic field there is strong and enhances the synchrotron radiation of the electrons and protons. Also the matter and photon densities are large enough in the core region, providing dense targets for proton-proton and proton-photon collisions and SSC scattering. In the energy range 10MeV − 40GeV, the emission is dominated by the neutral pion decay resulting from \( p − p \) inelastic collisions. Also, the resulting interactions between accelerated protons and scattered photons from synchrotron emission produce neutral pions and the gamma ray emission from these pion decays results in the tail seen in the SED for energies \( \geq 1 \)TeV.

6 DISCUSSION AND CONCLUSIONS

The multi-wavelength detection, from radio to gamma-rays, of non-thermal energy from galactic black hole binaries (BHBs or \( \mu \)QSRs) is clear evidence of an efficient production of relativistic particles and makes these sources excellent nearby laboratories to investigate and review particle acceleration theory in the surrounds of BH sources in general. Based on recent studies (GL05, [de Gouveia Dal Pino et al. 2010a,b], [Kadowaki & de Gouveia Dal Pino 2014]), we investigated here the role of magnetic reconnection in accelerating particles in the innermost regions of these sources, applying this acceleration model to reconstruct the spectral energy distribution (SED) of the BHBs Cyg X-1 and Cyg X-3.

According to GL05, particles can be accelerated to relativistic velocities in the surrounds of the BH, near the jet

basis, by a first-order Fermi process occurring in the magnetic reconnection discontinuity that forms where the magnetic field lines arising from the accretion disk encounter those anchored into the BH (Figure 1). This process becomes very efficient when these two magnetic line fluxes are squeezed together by enhanced disk accretion. This particle acceleration mechanism has been also tested successfully by means of numerical simulations (e.g. [Kowal, de Gouveia Dal Pino & Kowal 2013] and [del Valle et al. 2014]). This can be compared with the typical estimated acceleration time scale in diffusive shocks for the same environment conditions \( t_{\text{accel, shock}} \propto E \) (see eq. 5). We find a larger efficiency of the first mechanism in regions where magnetic discontinuities are dominant.

In fact, considering all the relevant leptonic and hadronic radiative loss mechanisms due to the interactions of the accelerated particles with the surrounding matter, magnetic and radiation fields in the core regions of Cyg X-1 and Cyg X-3, we compared the time scales of these losses with the acceleration time scales above and found larger energy cut-offs for particles being accelerated by magnetic reconnection than by a diffusive shock (see Figures 2 and 3 for Cyg X-1, and Figures 5 and 6 for Cyg X-3). These cut-offs have an important role in the determination of the energy distribution of the accelerated particles and therefore, in the resulting SED, and stress the potential importance of magnetic reconnection as an acceleration mechanism in the core regions of BHBs and compact sources in general.

In most astrophysical systems, synchrotron is known as a dominant mechanism to cool the electrons and for the sources studied here, its cooling rate is also larger than that of the other loss mechanisms in all electron energy range. In Cyg X-1, electrons gain energy up to 2 GeV (Figure 2), while the maximum energy that the electrons may reach in Cyg X-3 is around 3 GeV (Figure 5). In both cases, these values are larger than the possible ones obtained with shock acceleration.

Also, for both microquasars we find that \( p − \gamma \) is the dominant mechanism to cool the accelerated relativistic protons in most of the investigated energy range. Only for energies below \( \sim 2 \) TeV, the \( pp \) inelastic collisions are more efficient in the case of Cyg X-3. The calculated energy cut-off for protons obtained from the comparison of the \( p − \gamma \) cooling time with the magnetic reconnection acceleration time is 0.1PeV and 0.3PeV, for Cyg X-1 and Cyg X-3, respectively. In these \( p − \gamma \) processes, the synchrotron radiation is the dominant target photon field that interacts with the energetic protons, this because the magnetic field in the core region of these sources is relatively large, as calculated from Eq. 14.

We note that the maximum energy of the accelerated particles is also constrained by the size of the acceleration region, i.e., the particle Larmor radius, \( r_L = E/|e|B \), cannot be larger than the length scale of the acceleration zone. Considering the parameters employed in our model for both sources and \( \Delta R_X \) as the length scale of the acceleration zone, we find that the maximum energy to which the protons (and electrons) can be accelerated by magnetic reconnection is \( \sim 9 \) PeV, which is larger than the cut-off values.
obtained above. This value also reinforces the efficiency of this acceleration process.

We have also shown that, under fiducial conditions, the acceleration model developed here is capable of explaining the multi-wavelength non-thermal SED of both microquasars Cyg X-1 and Cyg X-3. The radio emission may result from synchrotron process in both cases.

The observed soft gamma-rays from Cyg X-1 are due to synchrotron and IC processes. The target photons for the IC come mainly from synchrotron emission (SSC). Neutral Pion decay resulting from $pp$ inelastic collisions may produce the high energy gamma rays in both systems, while the very high energy (VHE) gamma rays are the result of neutral pion decay due to photo-meson production ($p - \gamma$) in the core of these sources.

The importance of the $\gamma - \gamma$ absorption due to interactions with the photon field of the companion star for electron-positron pair production has been also addressed in our calculations. According to our results, the observed gamma-ray emission in Cyg X-1 in the range $3 \times 10^{10} - 3 \times 10^{11}$ eV cannot be produced in its core region (see also Romero, del Valle & Orellana 2010). In the case of Cyg X-3, we have found that the emission in the range of $50 \text{GeV} - 0.4 \text{TeV}$ is also fully absorbed in the core region by the same process. This suggests that in both sources, this emission is produced outside the core, probably along the jet, since at larger distances from the core the gamma ray absorption by the stellar companion decreases substantially. In fact, this is what was verified by Zhang, Xu & Lu (2014) in the case of Cyg X-1.

Other authors have proposed alternative scenarios to the one discussed here. The models of Piano et al. (2012), for instance, which were based on particle acceleration near the compact object and on propagation along the jet, indicate that the observed gamma-ray $\lesssim 10$ GeV in Cyg X-3 could be produced via leptonic (inverse Compton) and hadronic processes ($pp$ interactions). However, they have no quantitative estimates for the origin of the VHE gamma-ray upper limits at $\gtrsim 0.1$ TeV obtained by MAGIC. Sahakyan et al. (2013), on the other hand, assumed that the jet of Cyg X-3 could accelerate both leptons and hadrons to high energies and the accelerated protons escaping from the jet would interact with the hadronic matter of the companion star producing $\gamma-$rays and neutrinos. However, their model does not provide proper fitting in the TeV range either.

In the case of Cyg X-1, Zhang, Xu & Lu (2014) have employed a leptonic model to interpret recent Fermi LAT measurements also as due to synchrotron emission but produced along the jet and to Comptonization of photons of the stellar companion. The TeV emission in their model is attributed to interactions between relativistic electrons and stellar photons via inverse Compton scattering. According to them this process could also explain the MAGIC upper limits in the range of $30 \text{GeV} - 0.3 \text{TeV}$, i.e., the band gap in Figure 4. However, unlike the present work where we obtained a reasonable match due to $p - \gamma$ interactions, their model is unable to explain the observed upper limits by MAGIC in the very high energy gamma-ray tail.

Also with regard to Cyg X-1, we should note that the detection of strong polarized signals in the high-energy range of $0.4-2$ MeV by Laurent et al. (2011) and Jourdain et al. (2012) suggests that the optically thin synchrotron emission of relativistic electrons from the jet may produce soft gamma-rays. There are indeed some theoretical models that explain the emission in this range by using a jet model (Zdziarski et al. 2012; Malyshev et al. 2013; Zdziarski et al. 2014; Zhang, Xu & Lu 2014). Nevertheless, contrary to this view, Romero, Vieyro & Chaty (2014) argue that the MeV polarized tail may be originated in the coronal region of the core without requiring the jet. This study is therefore, consistent with the present model as it supports the coronal nuclear region for the origin of the non-thermal emission.

The results above clearly stress the current uncertainties regarding the region where the HE and VHE emission are produced in these compact sources. This work has tried to shed some light on this debate focussing on a core model with a magnetically dominated environment surrounding the BH, but a definite answer to this question should be given by much higher resolution and sensitivity observations which may be achieved in near future with the forthcoming Cherenkov Telescope Array (CTA) (Actis et al. 2011; Acharya et al. 2013; Sol et al. 2013).

We should also stress that there are two possible interpretations for the lack of clear evidence of detectable TeV emission in Cyg X-1 and Cyg X-3. On one hand, there may be a strong absorption of these photons by the ultraviolet (UV) radiation of the companion star (through the photon-photon process). On the other hand, the lack of emission may be due to the limited time of observation (Sahakyan et al. 2013). In our model, we verified that neutral pion decays due to $p - \gamma$ interactions at the emission region close enough to the central black hole, near the jet basis, could produce TeV gamma-rays. Because of the high magnetic field near the black hole, a large density synchrotron radiation field produced there could be a target photon field for the photo-meson production. These results predict that a long enough observation time and higher sensitivity would allow to capture substantial TeV $\gamma$-ray emission from these microquasars. This may be also probed by the CTA.

A final remark is in order. To derive the SEDs of the sources investigated here, we have assumed a nearly steady-state accelerated particle energy distribution at the emission zone. This assumption is valid as long as acceleration by magnetic reconnection is sustained in the inner disk region. Or in other words, as long as the disk accretion rate is large enough to push the magnetic field lines rising from the accretion disk to the magnetic field lines anchored into the BH. In microquasars, this should last no longer than the time the system remains in the outburst state, normally ranging from less than one day to several weeks.

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where $V_{\text{acc}}$ is the volume of the acceleration region and $L_{(e,p)}$ is the fraction of the magnetic reconnection power that accelerates the electrons and protons (see eq. 1 in the text). The injection particle spectrum is modified in the emission region due to energy losses. We assume that the minimum energy of the particles is given by $mc^2$, where $m$ is the rest mass of the particle. The maximum energy that the primary particles can attain is fixed by the balance of acceleration and the energy losses. Particles can gain energy up to a certain value $E_{\text{max}}$ for which the total cooling rate equals the acceleration rate.

The kinetic equation that describes the general evolution of the particle energy distribution $N(E,t)$ is the Fokker-Planck differential equation (Ginzburg & Syrovatskii 1995). We here use a simplified form of this equation. We employ the one-zone approximation to find the particle distribution, assuming that the acceleration region is spatially thin enough, so that we can ignore spatial derivatives in the transport equation. Physically, this means that we are neglecting the contributions to $N(E)$ coming from other regions than the magnetic reconnection region in the inner accretion disk zone in the surroundings of the BH.

We consider a steady-state particle distribution which can be obtained by setting $\frac{dN}{dt} = 0$ in the Fokker-Planck differential equation, so that the particle distribution equation is

$$N(E) = \left| \frac{dE}{dt} \right|^{-1} \int_{E}^{\infty} Q(E) dE. \quad (A3)$$

Here $-\frac{dE}{dt} \equiv E\frac{d\Gamma_{\text{cool}}}{dE}$. It is very interesting to note that if the energy losses are proportional to the particle energy ($\frac{dE}{dt} \propto E$), $N(E)$ does not change the injection spectrum and $N(E) \propto E^{-p}$, as in the $pp$ inelastic collisions or Bremsstrahlung cool processes. In such loss mechanisms like synchrotron and IC scattering, in the Thomson regime, the $N(E)$ is steeper because in these cases $\frac{dE}{dt} \propto E^{2}$ and $N(E) \propto E^{-(p+1)}$.

The spectrum would be harder if $dE/dt$ were constant as for ionization losses, $N(E) \propto E^{-(p-1)}$. In the case of IC scattering in the Klein-Nishina limit, $\frac{dE}{dt} \propto E^{-3}$ and so, the spectrum is even harder and $N(E) \propto E^{-(p-2)}$.

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