High-energy radiation from collisions of high-velocity clouds and the Galactic disc

Maria V. del Valle,1,2★ A. L. Müller1,3,4 and G. E. Romero1,5

1 Instituto Argentino de Radioastronomía (CONICET; CICPBA), C.C. No. 5, 1894 Villa Elisa, Argentina
2 Institute of Physics and Astronomy, University of Potsdam, D-14476 Potsdam, Germany
3 Institut für Kernphysik (IKP), Karlsruhe Institut of Technology, Karlsruhe, Germany
4 Instituto de Tecnologías en Detección y Astropartículas (CNEA, CONICET, UNSAM), Buenos Aires, Argentina
5 Facultad de Ciencias Astronómicas y Geofísicas, Universidad Nacional de La Plata, Paseo del Bosque s/n, 1900, La Plata, Argentina

Accepted 2017 November 16. Received 2017 November 16; in original form 2017 September 4

ABSTRACT
High-velocity clouds (HVCs) are interstellar clouds of atomic hydrogen that do not follow normal Galactic rotation and have velocities of a several hundred kilometres per second. A considerable number of these clouds are falling down towards the Galactic disc. HVCs form large and massive complexes, so if they collide with the disc a great amount of energy would be released into the interstellar medium. The cloud–disc interaction produces two shocks: one propagates through the cloud and the other through the disc. The properties of these shocks depend mainly on the cloud velocity and the disc–cloud density ratio. In this work, we study the conditions necessary for these shocks to accelerate particles by diffusive shock acceleration and we study the non-thermal radiation that is produced. We analyse particle acceleration in both the cloud and disc shocks. Solving a time-dependent two-dimensional transport equation for both relativistic electrons and protons, we obtain particle distributions and non-thermal spectral energy distributions. In a shocked cloud, significant synchrotron radio emission is produced along with soft gamma rays. In the case of acceleration in the shocked disc, the non-thermal radiation is stronger; the gamma rays, of leptonic origin, might be detectable with current instruments. A large number of protons are injected into the Galactic interstellar medium, and locally exceed the cosmic ray background. We conclude that under adequate conditions the contribution from HVC–disc collisions to the galactic population of relativistic particles and the associated extended non-thermal radiation might be important.

Key words: radiation mechanisms: non-thermal – ISM: clouds – cosmic rays.

1 INTRODUCTION
High-velocity clouds (HVCs) are a component of the neutral interstellar medium (ISM). These clouds are formed mainly by hydrogen and move with anomalous velocities (higher than the Galactic rotation velocity), with deviation velocities \( V_{\text{dev}} > 90 \text{ km s}^{-1} \).1 There are at least three suggested origins for these clouds: (i) material heated and transported to the Galactic halo by supernova explosions (such material cools and falls back to the Galactic plane) – this hypothesis is called the Galactic fountain; (ii) gas streams produced by tidal forces on nearby dwarf galaxies (e.g. the Magellanic Stream); (iii) low-metallicity matter of intergalactic origin that falls on to the Galaxy. The third, low-metallicity, component is thought to inject fresh material into the Galaxy for star formation and its existence is important to current models of Galactic evolution (e.g. Wakker & van Woerden 2013). In any case, HVCs are essential for understanding the flows of energy and mass towards and within the Galaxy. In the past few decades, their importance has grown considerably for two research fields: Galactic star formation and dark matter. The second field of research comes from the idea that at least some HVCs owe their existence to the presence of dark matter haloes (Blitz et al. 1999; Quilis & Moore 2001).

A large fraction of HVCs have negative (approaching) velocities, so they will reach the Galactic plane at some time, colliding with the disc. The impact of these clouds with the gas in the disc should release a large amount of energy into the ISM, between \( 10^{47} \) and \( 10^{52} \) erg. Such energetic events can trigger star formation episodes (Tenorio-Tagle 1981). For example, the supermassive HVC called the Smith cloud is thought to have collided with the Galactic disc approximately 70 Myr ago. It should cross the plane again in about...
27 Myr (Lockman et al. 2008). Numerical simulations of this collision suggest that, in order to have survived the impact, the Smith cloud should be embedded into a dark matter mini-halo (Nichols et al. 2014). Galyardt & Shelton (2016) studied the collision of this cloud and inferred the properties of putative dark matter. High- and intermediate-velocity clouds can also be used to trace cosmic rays (CRs) in the halo of the Milky Way (Tibaldo et al. 2015).

Several aspects of cloud–disc collisions as well as interactions of HVCs with their environment have been studied (e.g., see Chapter 12 in Waker & van Woerden 2013). However, the possibility of particle acceleration in these interactions has yet to be explored. Hedrick & Cox (1977) investigated the energy requirements for CR acceleration in the inflow of HVCs material towards the Galactic plane. Collisions of HVCs with the Galactic disc were studied as potential CR sources by Morfill & Tenorio-Tagle (1983), see also Romero & Paredes (2011). Blom et al. (1997) reported the first possible observational hint for extended MeV emission that may be associated with HVCs. In the current work, we explore the particle acceleration that might take place in the shocks formed in the collisions and we calculate the non-thermal radiative output. Preliminary results of this research have already been presented by Müller, Romero & del Valle (2017).

In the next section, we briefly discuss the properties of the shocks created by the clouds when they collide with the disc and we analyse the efficiency of diffusive shock acceleration (DSA) in these shocks. In Section 3, we present our model for the shock propagating through the cloud; we present and discuss the results of our calculations for this shock in Section 4. In Section 5, we describe the model for the shock, moving through the disc. The results for the shocked disc scenario are presented in Section 6. Finally, in Section 7 we give a short summary and offer our conclusions.

2 CLOUD–DISC COLLISIONS

HVCs have velocities in the range 100–500 km s$^{-1}$ and typical densities between 0.1 and 1.0 cm$^{-3}$. Even though the clouds are detected through the neutral H line, not all the gas is in neutral form. Ionized hydrogen is expected to constitute a large fraction of the material. Cloud radii vary greatly, from several pc to a few kpc, with a substantial number of clouds having radii greater than 50 pc. HVCs are thought to form large complexes. Metallicity and distance estimates are obtained through spectroscopy; the distance is ratioed by CRs in the halo of the Milky Way (Tibaldo et al. 2015).

Concerning the metallicity, most clouds present subsolar abundance estimates are obtained through spectroscopy; the distance is calculated for this shock in Section 4. In Section 5, we describe the model for the shock, moving through the disc. The results for the shocked disc scenario are presented in Section 6. Finally, in Section 7 we give a short summary and offer our conclusions.

2.1 Particle acceleration

The first-order Fermi mechanism$^2$ or diffusive shock acceleration is known to operate efficiently in very fast shocks, with velocities of the order of $10^{-2} c$ or higher. Recently, observations and numerical experiments have suggested that the first-order Fermi mechanism can also operate in slower shocks, with velocities $\sim 10^{-3} c$ (e.g. Caprioli & Spitkovsky 2014; Lee et al. 2015; Metzger et al. 2015). This seems to be supported by the detection of synchrotron radiation from young stellar objects (Carrasco-González et al. 2010; Rodríguez-Kamenetzky et al. 2016). However, there are some limitations on the DSA mechanism in slow shocks, as discussed below.

Shocks can be radiative or adiabatic, depending on the efficiency of the gas to lose energy through thermal radiation. In radiative shocks, the shocked material rapidly cools down, resulting in large compression factors. A very dense gas layer forms and the shock promptly slows down (e.g. Drake 2005). The efficiency $\eta$ to accelerate particles depends strongly on the shock velocity $V_s$, $\eta \propto (V_s/c)^2$, and hence this efficiency decays very fast in these shocks (although DSA in slightly radiative shocks might still operate). Moreover, if the post-shock density attains very high values, collisions and ionization losses might be catastrophic for particle acceleration (see below). However, an adiabatic shock propagates with approximately constant velocity through large spatial scales and its particle acceleration efficiency remains more or less constant. Because of this, we consider only adiabatic shocks in what follows.

In order to determine the nature of a shock, we compare the characteristic time-scale $t_{\text{chat}}$ of the physical processes involved with the time-scale of thermal losses $t_{\text{rad}}$. If $t_{\text{chat}} \ll t_{\text{rad}}$, then the shock is adiabatic. The time-scale $t_{\text{rad}}$ can be calculated as

$$t_{\text{rad}} = \frac{5}{3} \frac{P}{L},$$

where $P$ is the post-shock gas pressure, $T$ is the post-shock temperature, $T = 2 \times 10^{-9} V_s^2$ K (with $V_s$ in cm s$^{-1}$), $L = n^2 \Lambda(T)$, $n$ is the number density and $\Lambda(T)$ is the cooling function that can be fitted as a power law in $T$ (e.g. Myasnikov, Zheklov & Belov 1998)$^3$.

In the absence of significant magnetic field pressure, shocks cannot accelerate particles if their Mach number $M = V_s/C_s \leq \sqrt{3}$ and $M \leq 6$ for fully relativistic particles (Vink & Yamazaki 2014), where $C_s$ is the sound speed. The shock velocity should then be higher than $6C_s$. Even in the case of slow shocks, this constraint is not a limitation in this study because in the regions we are dealing with the temperatures are relatively low and the sound speed is much smaller than $100 \text{km s}^{-1}$.

When the ambient gas is very dense or the shock velocity is not high enough, ionization and Coulomb losses of low-energy particles can be catastrophic, halting the particle acceleration process. If effective, the acceleration has to be fast enough at supra-thermal energies to compete with the collisional losses. The losses will not suppress the acceleration at any energy if (O’C Drury, Duffy &
Here, adiabatic shocks, so we focus on clouds with high velocities—forward shocks (i.e. models CI and CII) and one case in which a strong shock (propagating through the disc), but this shock is radiative with the parameters adopted in models CI and CII. We then ignore the forward shocks in these models.

Magnetic fields can play an important role in cloud dynamics and collisions. There has only been one detection so far of magnetic fields in this type of object that establishes a lower limit to the field on the line of sight of $B_{\text{lim}} = 8 \mu G$ (Hill et al. 2013). We assume $B = 10 \mu G$ for both cloud models.

Because the clouds have relatively low densities, the condition given by equation (4) is widely fulfilled. We adopt a typical value for the cloud temperature of $T_c = 10^4 K$ (Wakker & van Woerden 2013).

It is not simple to calculate the maximum energy that particles achieve because DSA is a non-linear process; we can obtain an estimate by computing the balance between the energy gain and energy-loss rates. The relevant non-thermal radiative losses are synchrotron, relativistic bremsstrahlung and inverse Compton (IC) scattering with the cosmic background radiation (CMB) at $T = 2.7 K$, for electrons, and $p-p$ inelastic collisions for protons. Particles might also escape from the acceleration region by propagating through the cloud. The process occurs during the characteristic crossing time: $t_{\text{char}} = 2R_c/V_s$. Fig. 1 illustrates this scenario.

In Fig. 2, we show the ratio $t_{\text{rad}}/t_{\text{char}}$ in logarithmic scale as a function of density and shock velocity (see equation 3); the regions where $t_{\text{rad}}/t_{\text{char}} \geq 10$ are shown in red. These are the regions where the radiative losses are highly inefficient. All three models lie in the adiabatic zone, so all shocks under consideration are adiabatic. The collision of the cloud with the disc also creates a forward shock (propagating through the disc), but this shock is radiative with the parameters adopted in models CI and CII. We then ignore the forward shocks in these models.

**3 MODEL: SHOCKED HVC**

We model the cloud as an homogeneous sphere of radius $R_c = 10 \text{ pc}$. We assume that protons and electrons are accelerated by the shock propagating through the cloud. The process occurs during the characteristic crossing time: $t_{\text{char}} = 2R_c/V_s$. Fig. 1 illustrates this scenario.

| Model | $n_s$ ($\text{cm}^{-3}$) | $n_d$ ($\text{cm}^{-3}$) | $V_s$ ($\text{km s}^{-1}$) | $\mathcal{M}$ |
|-------|-----------------|-----------------|-----------------|----------------|
| CI    | 0.1             | 1.0             | 500             | 43             |
| CII   | 0.5             | 0.1             | 200             | 16             |
| D     | 1.0             | 1.0             | 500             | 50             |

Kirk 1996)

$$\left(\frac{V_s}{10^3 \text{ km s}^{-1}}\right)^2 \left(\frac{B}{1 \mu G}\right) \left(\frac{n}{1 \text{ cm}^{-3}}\right)^{-1} \geq 10^{-4} \max \left[\frac{1}{2}, (1 - \chi)\right].$$

(4)

Here, $n$ is the number density, $B$ is the magnetic field, $T_d$ is the gas temperature in units of $10^4 K$ and $\chi$ is the ionization fraction.

Another limitation for DSA in cold media is due to ion-neutral friction. A cold medium is not fully ionized and hence there is a large number of neutral atoms and molecules in the gas; the ion-neutral friction damps the turbulent inhomogeneities with which particles scatter during the acceleration process. Depending on the degree of ionization, these interactions can limit particle acceleration. Ion-neutral wave damping places no restriction on shock acceleration if the following condition is satisfied upstream (O’C Drury et al. 1996):

$$\left(\frac{V_s}{10^3 \text{ km s}^{-1}}\right)^3 \times 8 \times 10^{-3} \left(\frac{B}{1 \mu G}\right)^2 \times \left(\frac{n_s}{1 \text{ cm}^{-3}}\right) \left(\frac{n_i}{1 \text{ cm}^{-3}}\right)^{-2}.$$  

(5)

Here, $n_s$ and $n_i$ are the neutral and ion number density, respectively. If the latter condition is not fulfilled, particles would accelerate but only until they reach a break momentum. The maximum energy, in units of particle rest mass energy $m_c$, is given approximately by (Malkov, Diamond & Sagdeev 2011)

$$E_{\text{max}} / m_c \sim 10 \left(\frac{B}{1 \mu G}\right)^2 T_d^{-0.4} \times \left(\frac{n_s}{1 \text{ cm}^{-3}}\right)^{-1} \left(\frac{n_i}{1 \text{ cm}^{-3}}\right)^{-1/2}. $$

(6)

2.2 Models

We consider two cases of shocks moving through the cloud (reverse shocks) (i.e. models CI and CII) and one case in which a strong shock, induced by a HVC collision, propagates through the disc (a forward shock) (i.e. model D). We are interested in studying strong adiabatic shocks, so we focus on clouds with high velocities $\sim 500 \text{ km s}^{-1}$, as $V_s \propto V_c$ (see equations 1 and 2). Table 1 shows the main parameters of the models.

In the following sections, we describe the models and we analyse the properties of the shocks. First, we deal with the shocks in the cloud (models CI and CII) and then we discuss the shock propagating through the disc (model D).

**Table 1. Model parameters.**

**Figure 1.** Schematic diagram of the physical scenario considered for the HVC (not to scale).

**Figure 2.** The $t_{\text{rad}}/t_{\text{char}}$ ratio as a function of density and shock velocity. The red region corresponds to $t_{\text{rad}}/t_{\text{char}} \geq 10$ (i.e. adiabatic shocks).
diffusion; we adopt Bohm diffusion close to the shock. The values of $E_{\text{max-loss}}$ obtained for electrons and protons are shown in Table 2. Our quantitative estimates indicate that high energies are achievable (above 1 TeV).

The ionization degree of a cloud is not easy to estimate (e.g. van Woerden et al. 2004) and HVCs can vary from 99 per cent neutral to almost fully ionized. Sometimes a value $\chi_1 = 0.5$ is adopted. In other cases, however, there seems to be more ionized than neutral gas (Wakker et al. 2008). In the case of incomplete ionized clouds, ion-neutral friction might impose the maximum particle energy. In order to illustrate this last case, we take here a representative ionization degree of $\chi_1 = 0.5$. With this value, condition (5) is not fulfilled and the maximum energies allowed are obtained using equation (6). Because $E_{\text{max-frc}}$ results lower than $E_{\text{max-loss}}$ (see Table 2), ion-neutral friction halts the acceleration at the highest energies. Therefore, adopt $E_{\text{max}} = E_{\text{max-frc}}$ as the more realistic estimate.

### 3.1 Relativistic particle transport and emission

The transport of relativistic protons and electrons is supposed to occur in the test-particle approximation. The spectral energy distribution $N_\gamma$ of the particles obeys the following equation:

$$\frac{\partial N_\gamma}{\partial t} = D(E) \left\{ \frac{\partial}{\partial R} \left( R^2 \frac{\partial N_\gamma}{\partial R} \right) + \frac{1}{R^2 \sin \theta} \frac{\partial}{\partial \theta} \left( \sin \theta \frac{\partial N_\gamma}{\partial \theta} \right) \right\} - \frac{\partial}{\partial E} \int [P(R, \theta, E) N_\gamma] + Q_\gamma(R, \theta, E, t).$$

Here, $D(E)$ is the diffusion coefficient of the particles. $P(r, E)$ is the total radiative energy loss rate and $Q_\gamma(r, E, t)$ is the injection function. Given the geometry of the problem, we use a spherical coordinate system ($R, \theta, \phi$), with its origin at the cloud centre. The particle density function, $N_\gamma$, depends spatially only on $R$ and $\theta$, that is, $N_\gamma = N_\gamma(R, \theta, E, t)$.

Relativistic particles, accelerated at the shock, are injected in a surface $S_{\text{inj}} = 2\pi r^2$, that is moving with velocity $V_{\text{inj}} = V_c$ (see Fig. 1). We consider that the particles have a power-law distribution in energy of index $\alpha = 2$, as expected from DSA in strong non-relativistic shocks. The injection function is normalized according to the power available in relativistic particles $L_{\text{par}}$.

The total kinetic power of the shock is estimated as $L_{\text{kin}} = (1/2)\rho V_c^2 \Omega_{\text{sh}}$, where $\Omega_{\text{sh}}$ is the volume of the cloud. The shocks in models $CI$ and $CIH$ have Mach numbers $M \geq 10$, well above the limit of 6 (see Table 2). Numerical experiments show that non-relativistic shocks with such values of $M$ can transform 10 per cent or more of their kinetic power into relativistic particles through DSA (see Caprioli & Spitkovsky 2014), which is in agreement with other estimates (e.g. Ellison, Moebius & Paschmann 1990). Here, we take $L_{\text{par}} = 0.1 L_{\text{kin}}$, with $L_{\text{par}}$ equally divided between electrons and protons.

Beyond some spatial scale, the particle spatial diffusion changes from the Bohm regime in the acceleration region to a faster one (i.e. the diffusion coefficient increases). The acceleration process occurs within a region of linear size $L_{\text{acc}} = D_{\text{shock}}/V_c$. For the maximum energies considered here, $L_{\text{acc}} \leq 1$ pc. Because we are modelling a cloud of 10 pc, we are considering phenomena occurring on scales 10 times larger. Consequently, we expect the transition to a faster diffuse regime to occur within the cloud. We adopt a diffusion coefficient

$$D(E) = 10^{26} \left( \frac{E}{10 \text{ GeV}} \right)^{0.5} \text{ cm}^2 \text{s}^{-1},$$

similar to that of the ISM (e.g. Berezinskii et al. 1990) but slightly smaller, because the magnetic field in the HVCs is greater than in the ISM.

We solve equation (7) in a discrete grid $(E, R, \theta) \in [1 \text{ MeV}, 100 \text{ TeV}] \times [0, 10 \text{ pc}] \times [0, \pi]$, using the finite-volume method. The energy grid is logarithmically spaced, whereas the radial and polar grids are uniformly spaced. We use a grid resolution $(L, M, K) = (64, 32, 32)$. We integrate during $t_{\text{inj}} \equiv t_{\text{shock}}$. For a further description of the code, see the Appendix.

The particle distributions are interpolated into a three-dimensional spatial grid. We calculate the non-thermal radiation produced by the particles as they diffuse through the cloud. We neglect the increase of the magnetic field or density due to compression by the shock. This latter simplification produces an underestimation of at most a factor of 4. The results and their descriptions are presented in the next section.

### 4 RESULTS: SHOCKED HVC

The solution of the particle density distribution for protons of fixed energy $E = 10$ GeV is presented in Fig. 3. This figure shows the projected three-dimensional proton distributions in a two-dimensional map, as a function of the injection time. The injected protons suffer energy losses due to p–p inelastic collisions. The particle diffusion is not very fast at this energy, but its effects can be appreciated, especially in the last snapshots, in the shocked borders of the cloud. On these borders, the number of particles decreases.

The spectral energy distributions (SEDs) as a function of time are shown in Fig. 4, for model $CI$. The non-thermal radio emission peaks at $E_{\text{ph}} \sim 1.5 \times 10^3 \text{ eV} (\equiv 380 \text{ GHz})$ with luminosities greater than $10^{31} \text{ erg s}^{-1}$. Between radio and soft gamma rays, no significant non-thermal radiation is generated. The emission by protons reaches a maximum value at $E_{\text{ph}} \sim 5 \times 10^8 \text{ eV}$. However, in this region of the SED, the contribution from the relativistic electrons dominates, with a peak at $E_{\text{ph}} \sim 10^{10} \text{ eV}$ and luminosities $\sim 5 \times 10^{30} \text{ erg s}^{-1}$. The part of the SEDs due to electrons keeps approximately the same energy losses due to p–p inelastic collisions. The particle diffusion

Four We assume that the diffusion coefficient depends only on the particle energy.
at \( E_{\text{ph}} > 1 \) TeV. The SED greater luminosities are not reached at \( t = t_{\text{inj}} \), but a little earlier.

The SEDs as a function of time for the model \( \text{CII} \) are shown in Fig. 5. In this case, the luminosities are lower because of the smaller power in relativistic particles. These SEDs achieve greater luminosities at the final integration time, \( t = t_{\text{inj}} \). The non-thermal radio emission peaks at lower energies, \( E_{\text{ph}} \sim 3 \times 10^{-5} \) eV (\( \equiv 0.3 \) GHz). The emission produced by protons reaches a maximum around the same energy as in the previous case. The gamma luminosities do not go beyond energies of 1 TeV and no steepening in the SEDs appears at high energies. Also, IC radiation dominates the spectrum at soft gamma rays, with a power in excess of \( 10^{30} \) erg s\(^{-1}\). At the highest energies, \( 10^{10} < E_{\text{ph}} < 10^{12} \) eV, the emission is greater than \( 10^{28} \) erg s\(^{-1}\).

Fig. 6 shows the time variation of the total emitted power for the four main non-thermal radiative processes: synchrotron, relativistic bremsstrahlung, IC scattering and p–p interactions. This figure corresponds to model \( \text{CI} \). The power grows slowly for the four mechanisms, varying by almost two orders of magnitude between the initial time and the time when the maximum is reached.
This maximum occurs almost at the same time for the leptonic processes, after the maximum of the p–p emission. Clearly, synchrotron radiation is the most efficient non-thermal mechanism here, followed by IC. The maximum power emitted is $~5.6 \times 10^{32}$ erg s$^{-1}$ for synchrotron. IC reaches $~3 \times 10^{33}$ erg s$^{-1}$, followed by relativistic bremsstrahlung that reaches $~2.5 \times 10^{30}$ erg s$^{-1}$. Finally, we have $~7 \times 10^{29}$ erg s$^{-1}$ for p–p.

In the case of model CII (not shown here), the maximum luminosities of bremsstrahlung and p–p are greater, of $~8.7 \times 10^{30}$ and $~2.9 \times 10^{30}$ erg s$^{-1}$, respectively. The total synchrotron and IC power are lower than in the case of model CI, being $~3.4 \times 10^{31}$ and $~10^{31}$ erg s$^{-1}$, respectively.

### 4.1 Discussion

HVCs form large complexes, with sizes 10 times, or more, the size of the individual clouds modelled here. A large cloud complex might fragment into smaller pieces during the collision process. In such a case, the collective emission of all impacts can be one order of magnitude greater than the values we obtained for a single cloud.

The non-thermal radio emission is the most significant radiative output produced in the collision. This emission peaks near 380 GHz in the examples investigated here, far from the observed radio emission of HVCs at 21 cm ($\sim 1.4$ GHz = $5.9 \times 10^{-6}$ eV). The collective luminosity can be as high as $~5 \times 10^{33}$ erg s$^{-1}$, which is potentially detectable considering distances of the order of kpc.

The SED at soft gamma rays peaks near $E_{\text{th}} \sim 1$ GeV, close to the sensitivity peak of the Large Area Telescope (LAT) of the gamma satellite Fermi. For a source at $d \sim 1$ kpc, on the Galactic plane, Fermi might detect sources at $E_{\text{th}} = 10^{3}$ MeV over $~5 \times 10^{33}$ erg s$^{-1}$, for the one calendar year all-sky survey. In the collective case, the luminosity might be detectable, and a longer integration time could result in a $5\sigma$ detection.

At higher energies, around $E_{\text{th}} \sim 1$ TeV, the system of Cherenkov telescopes MAGIC can detect a source with power over $~2.3 \times 10^{31}$ erg s$^{-1}$ at 1 kpc (see Aleksić et al. 2016). The future array of Cherenkov telescopes CTA, at this same $E_{\text{th}}$, would have a sensitivity one order of magnitude higher; hence, CTA might be able to detect the gamma rays produced in the shocked cloud for a single cloud–disc collision.

The column densities of the HVCs are not too high. If the column density to the source (i.e. a cloud in the Galactic plane at 1 kpc) is too high, then the non-thermal emission produced by the Galactic CRs (electrons and protons) might be higher than the total luminosity of the shocked cloud, which will then be hidden by the background noise.

The proton CR flux in the Galaxy is given by (e.g. Simpson 1983)

$$J_{\text{CR}}^p(E) = 2.2 \left( \frac{E}{\text{GeV}} \right)^{-2.75} \text{cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{GeV}^{-1}.$$  \hspace{1cm} (9)

For electrons, we consider that $J_{\text{CR}}^e(E) = J_{\text{CR}}^p(E)/100$, as indicated from observations (e.g. Berezhikskii et al. 1990). Using these fluxes, we computed the relativistic bremsstrahlung and p–p radiation expected from the background. For a column density of $~10^{19}$ cm$^{-2}$ (Dickey & Lockman 1990), both emissions lie orders of magnitude lower than the contributions from the cloud, during practically all the integration time. In the case of a denser column density, $~10^{20}$ cm$^{-2}$, only the background p–p emission is higher than that of the cloud; however, this happens only at low energies, between 10 MeV and 10 GeV (see the grey solid line in Fig. 4).

The non-thermal radiation is the most efficient non-thermal mechanism here, followed by IC. The maximum power emitted is $~3 \times 10^{33}$ erg s$^{-1}$, followed by relativistic bremsstrahlung that reaches $~10^{30}$ erg s$^{-1}$. Finally, we have $~7 \times 10^{29}$ erg s$^{-1}$ for p–p.

In the case of model CII (not shown here), the maximum luminosities of bremsstrahlung and p–p are greater, of $~8.7 \times 10^{30}$ and $~2.9 \times 10^{30}$ erg s$^{-1}$, respectively. The total synchrotron and IC power are lower than in the case of model CI, being $~3.4 \times 10^{31}$ and $~10^{31}$ erg s$^{-1}$, respectively.

4.1 Discussion

HVCs form large complexes, with sizes 10 times, or more, the size of the individual clouds modelled here. A large cloud complex might fragment into smaller pieces during the collision process. In such a case, the collective emission of all impacts can be one order of magnitude greater than the values we obtained for a single cloud.

The non-thermal radio emission is the most significant radiative output produced in the collision. This emission peaks near 380 GHz in the examples investigated here, far from the observed radio emission of HVCs at 21 cm ($\sim 1.4$ GHz = $5.9 \times 10^{-6}$ eV). The collective luminosity can be as high as $~5 \times 10^{33}$ erg s$^{-1}$, which is potentially detectable considering distances of the order of kpc.

The SED at soft gamma rays peaks near $E_{\text{th}} \sim 1$ GeV, close to the sensitivity peak of the Large Area Telescope (LAT) of the gamma satellite Fermi. For a source at $d \sim 1$ kpc, on the Galactic plane, Fermi might detect sources at $E_{\text{th}} = 10^{3}$ MeV over $~5 \times 10^{33}$ erg s$^{-1}$, for the one calendar year all-sky survey. In the collective case, the luminosity might be detectable, and a longer integration time could result in a $5\sigma$ detection.

At higher energies, around $E_{\text{th}} \sim 1$ TeV, the system of Cherenkov telescopes MAGIC can detect a source with power over $~2.3 \times 10^{31}$ erg s$^{-1}$ at 1 kpc (see Aleksić et al. 2016). The future array of Cherenkov telescopes CTA, at this same $E_{\text{th}}$, would have a sensitivity one order of magnitude higher; hence, CTA might be able to detect the gamma rays produced in the shocked cloud for a single cloud–disc collision.

The column densities of the HVCs are not too high. If the column density to the source (i.e. a cloud in the Galactic plane at 1 kpc) is too high, then the non-thermal emission produced by the Galactic CRs (electrons and protons) might be higher than the total luminosity of the shocked cloud, which will then be hidden by the background noise.

The proton CR flux in the Galaxy is given by (e.g. Simpson 1983)

$$J_{\text{CR}}^p(E) = 2.2 \left( \frac{E}{\text{GeV}} \right)^{-2.75} \text{cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{GeV}^{-1}.$$  \hspace{1cm} (9)

For electrons, we consider that $J_{\text{CR}}^e(E) = J_{\text{CR}}^p(E)/100$, as indicated from observations (e.g. Berezhikskii et al. 1990). Using these fluxes, we computed the relativistic bremsstrahlung and p–p radiation expected from the background. For a column density of $~10^{19}$ cm$^{-2}$ (Dickey & Lockman 1990), both emissions lie orders of magnitude lower than the contributions from the cloud, during practically all the integration time. In the case of a denser column density, $~10^{20}$ cm$^{-2}$, only the background p–p emission is higher than that of the cloud; however, this happens only at low energies, between 10 MeV and 10 GeV (see the grey solid line in Fig. 4).

From Fig. 2 it can be seen that almost all the parameter space considered lies in the adiabatic region (red). Thus, we can have adiabatic shocks for a number of different parameters. Here we investigate the case of clouds with $V_c = 500 \text{km s}^{-1}$ as an extreme case; the majority of HVCs might have lower velocities. We analyse qualitatively what would be expected when changing the density and the velocity of the cloud.

In our model, the cloud velocity and the density determine the shock velocity and the power in relativistic particles. This last dependence is linear with the cloud density. A change in the power in relativistic particles is directly proportional to the non-thermal luminosity produced. The gamma emission from p–p collisions is also linearly proportional to the density and will vary accordingly; however, this is not the dominant non-thermal radiative process in these sources. Given that the average number density of HVCs does not change in many orders of magnitude, for a fixed shock velocity the number of relativistic particles is not expected to vary greatly from source to source.

The velocity of the shock is a more sensitive parameter of our model, and it is proportional to the cloud velocity. The shock velocity enters in the power in relativistic particles as $\propto V_c^2$. For HVCs with velocities between 100 and 500 km s$^{-1}$, and fixed density, the power in relativistic particles varies by a factor of 25. This corresponds to a variation of a factor of 25 in the non-thermal emission. Also, the acceleration rate is $\propto V_c^2$, and a variation of a factor of 25 is also expected in $E_{\text{max-loss}}$. If the maximum energies are determined by the losses and not by the ion-neutral damping (a highly ionized cloud), then the maximum energies the particles can reach vary by one order of magnitude.

Here, we use an ionization factor of $\chi_i = 0.5$. For the parameters adopted in models CI and CII, in the case of a fully or almost fully ionized cloud, ion-neutral friction is not relevant, and the maximum energies are higher, given by $E_{\text{max-loss}}$. In an almost complete neutral cloud, the maximum energies particles attain would be slightly lower than those in Table 2 (see dependences in equation 6).

We study only adiabatic shocks because the DSA theory is well understood in this regime. Radiative shocks are ubiquitous in our Galaxy; these shocks might also accelerate particles via DSA. However, the radiation losses modify the shocks and it is very complex.
to study the acceleration mechanism operating in such a regime. It would be too speculative to assume DSA in radiative shocks without a careful analysis. However, the acceleration of particles and/or re-acceleration of pre-existing CRs in radiative shocks produced in HVC-disc collision systems cannot be ruled out.

5 MODEL: SHOCKED DISC

In order to model the effects of the collision of the disc, we assume that the forward shock, which propagates through an homogeneous disc, injects particles as a point source (see Fig. 7). The region of impact in the disc is modelled as a cylinder of radius \( r_d = 100 \) pc, height \( h_d = r_d = 100 \) pc and density \( n_d = 0.1 \) cm\(^{-3}\). The sound velocity in the warm ISM is of the order of \( 10 \) km s\(^{-1}\) (e.g. Draine \\& Lazarian 1998). The shock Mach number is then \( M \sim 50 \gg 6 \) (see Table 1). The characteristic crossing time is \( t_{\text{char}} = h_d / V_s \), so \( t_{\text{rad}} \gg t_{\text{char}} \) and the shock is adiabatic (see Section 2.1). We adopt a magnetic field \( B_d = 4 \) \( \mu \)G for the compressed medium.

The forward shock injects protons and electrons along the characteristic time \( t_{\text{char}} \). Fig. 8 shows a schematic diagram of the acceleration scenario. As before, we estimate the particles’ maximum energies by comparing the energy loss and gain rates. The relevant non-thermal radiative losses of electrons are due to synchrotron radiation, relativistic bremsstrahlung and IC scattering with the background radiation fields. The most relevant interstellar radiation fields are the CMB, the ambient infrared (IR) radiation field (mainly produced by thermal emission from interstellar dust grains) and the ultraviolet (UV) contribution from the integrated stellar radiation (e.g. Maciel 2013). Protons lose energy only through p–p inelastic collisions. Particles, as in the previous cases, might escape the acceleration region, of size \( l_{\text{acc,b}} \), by diffusion and now they might also be drawn away from the acceleration zone, advected by the material that flows through the sides of the cloud with a velocity \( V_{\text{adv}} \sim V_s / 4 \) (Fig. 8).

When a HVC approaches the disc, it will find a composition of the different ISM phases that is mostly neutral, but with a strongly varying (0–100 per cent) ionization fraction that is irregularly distributed on scales smaller than 100 pc. We consider the limiting cases: a fully ionized disc and a 99 per cent neutral disc (i.e. \( \chi_i \sim 0.01 \)).

Fig. 9 shows the losses and acceleration time-scales of electrons (upper panel) and protons (lower panel), for model D.

![Figure 7. Schematic diagram of the physical scenario adopted for the Galactic disc (not to scale). The gridded purple surface indicates the computational plane \( r, z \). The cloud–shock system is taken as a punctual injector.](image)

![Figure 8. Schematic diagram of the acceleration scenario for the shock propagating through the disc, model D (not to scale).](image)

![Figure 9. Energy losses and acceleration time-scales of electrons (upper panel) and protons (lower panel), for model D.](image)

Table 3. Estimates of parameters in model D.

| \( E_{\text{max-loss}} \) (GeV) | \( E_{\text{max-fric}} \) (GeV) | \( L_{\text{par}} \) (erg s\(^{-1}\)) | \( t_{\text{inj}} \) (Myr) |
|---|---|---|---|
| \( e \) \( 7 \times 10^3 \) | \( p \) \( 7 \times 10^3 \) | \( e \) \( 5 \times 10^1 \) | \( p \) \( 5 \times 10^4 \) | \( L_{\text{par}} \) \( 2 \times 10^{36} \) | \( t_{\text{inj}} \) \( 2 \times 10^{-1} \) |

We only consider the details of the acceleration region for estimating the maximum energies. In the calculations of the transport of relativistic particles and non-thermal emission, the source is treated as a punctual injector.
that $E_{\text{max-fir}}$ for protons is greater than $E_{\text{max-loss}}$, and then the proton maximum energies for a neutral or fully ionized disc are identical. For electrons, the maximum energy for $\chi_1 \sim 0.01$ is two orders of magnitude lower than for a fully ionized disc.

5.1 Relativistic particle transport and emission

A cylindrical coordinate system $(r, z)$ is the most natural system to describe particle transport in the disc. The $z$-axis is determined by the direction of the cloud velocity. The transport equation for relativistic protons and electrons, in cylindrical coordinates, is

$$\frac{\partial N_p}{\partial t} = D(E) \left[ \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial N_p}{\partial r} \right) + \frac{\partial^2 N_p}{\partial z^2} \right] - \frac{\partial}{\partial t} \left[ P(r, z, E, t) N_p \right] + Q_p(r, z, E, t),$$

(10)

where the terms are the same as those in equation (7). We solve the former equation in a $(r, z)$ plane (see Fig. 7) with $0 \leq z \leq z_{\text{max}} \equiv h_1$ and $0 \leq r \leq r_{\text{max}} \equiv r_d$. The energy grid is the same as in the previous cases. We use a grid resolution $$(64, 64, 64)$$. Again, we integrate during $t_{\text{inj}} \equiv t_{\text{char}}$ (see the Appendix).

The particles accelerated at the shock are injected in the disc, with a velocity $V_{\text{inj}} \equiv V_c$. We consider, as before, that the particles have a power-law distribution in energy of index $\alpha = 2$. Then, the injection term is $Q_p(r, z, E, t) = Q_{\text{norm}} E^{-\alpha} \delta(X - X_{\text{inj}})$, where $X_{\text{inj}}$ is the shock position at time $t$, that is, $X_{\text{inj}} = (0, V_{\text{inj}}t)$. $Q_{\text{norm}}$ is the normalization factor, which depends on the power available in relativistic particles $L_{\text{rad}}$. We do not take into account any physical details of the shock-cloud system as we model the source as a point-like injector.

A fraction of the kinetic energy of the cloud is transferred to the forward shock in the collision. This fraction is large, so we approximate it as 1/2. Then, the power in the shock results $L_{\text{kin}} = (1/4) \rho c V_c^2 \Omega_c / t_{\text{char}}$. As mentioned before, 10 per cent or more of the shock power goes into relativistic particles in the DSA process. Consequently, we adopt $L_{\text{rad}} = 0.1 L_{\text{kin}}$, with $L_{\text{rad}}$ equally divided between electrons and protons (see Table 3).

Beyond some spatial scale, the diffusion coefficient becomes higher than in the Bohm approximation. Because we are interested in the large-scale phenomena, we adopt a diffusion coefficient typical of the ISM:

$$D(E) = 10^{27} \left( \frac{E}{10 \text{ GeV}} \right)^{0.5} \text{ cm}^2 \text{ s}^{-1}.$$  

With the time-dependent particle distributions obtained, we calculate the non-thermal radiation produced by synchrotron, IC scattering, bremsstrahlung, and by neutral pion decay due to p–p inelastic collisions.

We are interested in comparing the locally injected protons with that from the background CR population. In order to obtain the background CR distribution consistently with the parameters adopted, we solve the transport equation in the steady state with a null injection function, with the initial condition $N_p^0(t = 0) = 4\pi/c J_{\text{CR}}^p(E)$ (given in equation 9) and with boundary conditions matching the proton flux of the Galactic CR protons $J_{\text{CR}}^p$.

6 RESULTS: SHOCKED DISC

Fig. 10 shows maps of the evolution of protons with $E = 1$ GeV (top) and $E = 1$ TeV (bottom) as they are injected into the disc. In order to produce these maps, we project the three-dimensional space $(r, z, \phi)$, with $0 \leq \phi \leq 2 \pi$ the azimuthal angle, into a $(x, y)$ Cartesian plane. As the particles are injected, they diffuse through the disc; the more energetic particles diffuse faster. During the integration time, most of these particles do not reach the boundaries of the region, so they are concentrated around the axis in a zone of radius $\sim 20$ pc for the protons with $E = 1$ GeV, and of $\sim 66$ pc for those with $E = 1$ TeV.

The corresponding non-thermal SEDs for the cases $\chi_1 \sim 1$ and $\chi_1 \sim 0.01$ are shown in Fig. 11, for three different times during the integration period. We only calculate the IC emission due to interactions with the cosmic background, because this is by far the dominant contribution (see Fig. 9). The emission reaches its maximum when the injector is located in the middle of the integration region. The contribution from protons, as expected, is identical in both cases. The main contributions come from the electrons, with a higher luminosity at radio wavelengths due to the synchrotron radiation that is of the order of $10^{33}$ erg s$^{-1}$ for $\chi_1 \sim 1$ and $10^{32}$ erg s$^{-1}$ for $\chi_1 \sim 0.01$. This difference in the luminosity is because in the neutral case there is a lack of energetic electrons. Also, the synchrotron emission reaches smaller energies in the neutral disc case, as expected from an electron population with lower maximum energy.

The gamma-ray emission is dominated by IC, reaching a maximum luminosity of $\sim 5 \times 10^{31}$ erg s$^{-1}$ at soft gamma rays in both cases, but in the case with $\chi_1 \sim 0.01$ this maximum is slightly shifted to lower energies. The contributions from relativistic bremsstrahlung and from p–p collisions are negligible.

In the middle panel of Fig. 11, in grey, we plot the Fermi sensitivity curve for a source at $d \sim 1$ kpc. The IC luminosity lies above the curve, even in the case of a single collision. Hence, the gamma radiation from an event of these characteristics might be detectable at energies around $\sim 10$ GeV for both $\chi_1 \sim 1$ and $\chi_1 \sim 0.01$. Some unidentified Fermi sources might correspond to this type of object.

In the case of multiple events, as discussed in the previous scenario, the luminosity can increase by one order of magnitude, making the event detectable on a wider energy range. Also, the collective emission might be detectable, at least around the energies of maximum Fermi sensitivity, at greater distances (i.e. up to $d \sim 5$ kpc).

The protons lose only a small fraction of their energy by p–p in the cloud, and they diffuse into the surroundings, adding their energy to the local CR sea. In what follows, we briefly discuss the implications of this.

6.1 Discussion

In order to compare locally the injected protons with those from the CR background, we calculate this latter contribution as explained in Section 5.1; the results are shown over the maps in Fig. 10. The white dashed curves indicate the region where the local proton flux is equal to the background CR protons, and hence in the region inside the curve the protons accelerated at the shocked disc exceed the background. This region, around the injection axis, is greater for later times, and it extends from the axis $\sim 10$ pc for $E = 1$ GeV up to $\sim 33$ pc in the case of higher energies.

The locally injected protons dominate over the background on a considerable region, especially at high energies. This is expected because of the softer distribution of CRs compared with the local particles. As a result, the shock propagation through the disc enhances the local number of energetic protons. This effect can be even stronger when considering collective effects (i.e. a bigger HVC fragmenting in smaller pieces of the size studied here; see Section 4.1). Such a case is illustrated by the golden (outer) dashed curves in Fig. 10, where we consider that the injected number of...
Figure 10. Evolution of the number of protons with $E = 1$ GeV (top) and $E = 1$ TeV. These maps show the projection of the number of particles in the cylindrical disc on to a $x$, $y$-plane with $y$ in the direction of the cloud’s motion. Time evolves from left to right.

Figure 11. SED for three different injection times for model D in the case of a fully ionized disc, $\chi_i \sim 1$ (top) and an almost neutral disc, $\chi_i \sim 0.01$ (bottom). The left panel shows the non-thermal SED for $t = 0.1 t_{\text{inj}}$, the middle panel shows the SED at $t = 0.5 t_{\text{inj}}$ and the right panel corresponds to the final integration time $t = 1.0 t_{\text{inj}}$. 
protons is one order of magnitude higher. The region where the local flux dominates over the background extends beyond 66 pc for protons of $E = 1 \text{ TeV}$.

The rate of mass injected in the Galaxy by HVCs impacts is $0.5 M_\odot \text{ yr}^{-1}$ (Richter 2012). The total power released, if all the clouds have the same velocity $V_c$, is

$$P_{\text{HVCs}} \sim 2 \times 10^{40} \left(\frac{V_c}{250 \text{ km s}^{-1}}\right)^2 \text{ erg s}^{-1}. \quad (12)$$

Assuming that half of this power is transfer in the collisions to disc-propagating shocks, we can estimate the power in relativistic particles. If the shocks accelerate particles with a 10 per cent efficiency, then for an average cloud velocity of $V_c \sim 250 \text{ km s}^{-1}$, the power in relativistic particles is $P_{\text{CR}} \sim 10^{39} \text{ erg s}^{-1}$. This power is non-negligible considering that the power in the Galaxy from CRs is $\sim 5 \times 10^{40} \text{ erg s}^{-1}$. We conclude that up to 10 per cent of the CRs in the Galaxy can be produced by impacts of HVCs.

7 SUMMARY AND CONCLUSIONS

In this work, we investigated the non-thermal effects of the collision of a HVC with the Galactic disc. We analysed the properties of the shocks produced in the interaction, concluding that under some general conditions DSA can operate efficiently. We considered two different scenarios for computing the non-thermal emission: a reverse shock propagating through the cloud and a forward shock propagating through the Galactic disc.

We solved the transport equation for electrons and protons in a spherical cloud, for two different sets of parameters. We found that significant non-thermal radio emission occurs, with maximum luminosities around $\sim 10^{35} \text{ erg s}^{-1}$; we also found a moderate soft gamma component, with maximum order of the power of $10^{36} \text{ erg s}^{-1}$. In the case of multiple impacts, the corresponding luminosities can be an order of magnitude higher.

For the case of an adiabatic shock propagating through the disc, we solved the transport equation of particles in cylindrical coordinates. The leptonic contributions dominate the SEDs, with synchrotron radiation being at the level of $\sim 10^{35} \text{ erg s}^{-1}$ and gamma-ray emission, from IC up-scattering of the cosmic background, at $\sim 10^{36} \text{ erg s}^{-1}$. We also compared the number of protons that are locally accelerated with the Galactic CR population, finding that the local component of protons (especially those of the highest energies) dominates over the background on a substantial region.

The impact of HVCs with the Galactic disc releases a great amount of energy into the ISM. We have shown that under some reasonable conditions a fraction of that energy can be converted into non-thermal radiation and energetic particles. Furthermore, under some circumstances the non-thermal emission might be detectable and the power in energetic particles might form a non-negligible contribution of the global population of Galactic CRs.

ACKNOWLEDGEMENTS

The authors would like to thank Dr Reinaldo Santos-Lima for his help with numerical affairs. MVdV acknowledges partial support from the Alexander von Humboldt Foundation. GER acknowledges support from the Argentine Agency CONICET (PIP 2014-00338) and the Spanish Ministerio de Economía y Competitividad (MINECO/FEDER, UE) under grants AYA2013-47447-C3-1-P and AYA2016-76012-C3-1-P.

REFERENCES

Aleskić J. et al., 2016, Astropart. Phys., 72, 76
Bell A. R., 1978, MNRAS, 182, 147
Berezinskii V. S., Bulanov S. V., Dogiel V. A., Ptuskin V. S., 1990, Astrophysics of Cosmic Rays. North-Holland, Amsterdam
Blitz L., Spergel D. N., Teuben P. J., Hartmann D., Burton W. B., 1999, ApJ, 514, 818
Blom J. J. et al., 1997, A&A, 321, 288
Caprioli D., Spitkovsky A., 2014, ApJ, 783, 91
Carrasco-González C., Rodríguez L. F., Anglada G., Martí J., Torrelles J. M., Osorio M., 2010, Sci. 330, 1209
del Valle M. V., Romero G. E., Santos-Lima R., 2015, MNRAS, 448, 207
Dickey J. M., Lockman F. J., 1990, ARA&A, 28, 215
Draine B. T., Lazarian A., 1998, ApJ, 494, L19
Drake R. P., 2005, Ap&SS, 298, 49
Ellison D. C., Moebius E., Paschmann G., 1990, ApJ, 352, 376
Galyardt J., Shelton R. L., 2016, ApJ, 816, L18
Hedrick D., Cox D. P., 1977, ApJ, 215, 208
Hill A. S., Mao S. A., Benjamin R. A., Lockman F. J., McClure-Griffiths N. M., 2013, ApJ, 777, 55
Lee H. M., Kang H., Ryu D., 1996, ApJ, 464, 131
Lee S.-H., Patnaude D. J., Raymond J. C., Nagataki S., Slane P. O., Ellison D. C., 2015, ApJ, 806, 71
Lockman F. J., Benjamin R. A., Heroux A. J., Langston G. I., 2008, ApJ, 679, L21
Maciel W. J., 2013, Astrophysics of the Interstellar Medium. Springer-Verlag, Berlin
Malkov M. A., Diamond P. H., Sagdeev R. Z., 2011, Nature Communications, 2, 194
Metzger B. D., Finzell T., V Dunn T., Hascoët R., Beloborodov A. M., Chomiuk L., 2015, MNRAS, 450, 2739
Morfill G. E., Tenorio-Tagle G., 1983, Space Sci. Rev., 36, 93
Müller A. L., Romero G. E., del Valle M. V., 2017, in Aharonian F. A., Hofmann W., Rieger F. M., eds, AIP Conf. Proc. Vol. 1792, High Energy Gamma-Ray Astronomy: 6th International Meeting on High Energy Gamma-Ray Astronomy. American Institute of Physics, New York, 040007
Myasnikov A. V., Zhekov S. A., Belov N. A., 1998, MNRAS, 298, 1021
Nichols M., Mirabal N., Agertz O., Lockman F. J., Bland-Hawthorn J., 2014, MNRAS, 442, 2883
O’C Drury L., Duffy P., Kirk J. G., 1996, A&A, 309, 1002
Quilis V., Moore B., 2001, ApJ, 555, L95
Richter P., 2012, ApJ, 750, 165
Rodríguez-Kamenetzky A., Carrasco-González C., Araudo A., Torrelles J. M., Anglada G., Martí J., Rodríguez L. F., Valotto C., 2016, ApJ, 818, 27
Romero G. E., Paredes J. P., 2011, Introducción a la astrofísica relativista, Publicaciones e Ediciones de la Universitat de Barcelona
Simpson J. A., 1983, Annual Review of Nuclear and Particle Science, 33, 1227
Tenorio-Tagle G., 1981, A&A, 94, 338
Tibaldo L. et al., 2015, ApJ, 807, 161
van Woerden H., Wakker B. P., 2004, ApJ, 750, 165
van Woerden H., Wakker B. P., 2013, in Oswalt T. D., Gilmore G., eds, Planets, Stars and Stellar Systems Vol. 5. Springer-Verlag, Berlin, p. 659
Wakker B. P., van Woerden H., 2013, in Oswalt T. D., Gilmore G., eds, Planets, Stars and Stellar Systems Vol. 5. Springer-Verlag, Berlin, p. 587
Wakker B. P., York D. G., Wilhelm R., Barentine J. C., Richter P., Beers T. C., Ivezić Z., Howk J. C., 2008, ApJ, 672, 298

APPENDIX: TRANSPORT OF ENERGETIC PARTICLES

In order to solve the transport of relativistic particles and to calculate the subsequent non-thermal emission, we have used a modular code
The transport equations (7) and (10) were solved in spherical and cylindrical coordinates, respectively, allowing a reduction of the dimensionality of the problem given the existing symmetries.

The transport equations for electrons and protons were evolved simultaneously using the finite-volume method. We adopted a discrete grid \((E, R, \theta)\) in spherical coordinates and \((E, r, z)\) in cylindrical coordinates, where \((R, \theta)/(r, z)\) are the usual spatial coordinates in the corresponding system and \(E\) is the energy of the particles. The energy grid was logarithmically spaced, whereas the spatial grids were sampled uniformly.

Initially, at time \(t = 0\), we considered no particles inside the domain. We imposed that there were no particles outside the energy bounds. These limits did not influence the system evolution, because the upper limit was above the maximum energy of the injected particles, at the same time that the advection in the energy space (the term of energy losses) is always towards smaller energies. The lower bound was physically fixed by the particles’ rest mass.

For the spatial boundary conditions, we considered zero particles outside the domain (i.e. for \(R > r_c\) in the spherical case, and for \(r > r_d\) and \(z > r_d\) in cylindrical coordinates). Because of the azimuthal symmetry in the case of spherical coordinates at \(\theta = 0\) and \(\theta = \pi\), there we imposed outflow boundary conditions. In the case of cylindrical coordinates, at the inner boundary condition for \(r\) we adopted axial symmetry.

The numerical integration was performed through the operator splitting method. Each time-step integration evolved the particle density distribution on the grid through three substeps: first we integrated the losses, then the spatial diffusion and finally we added the source term.

The loss term is an advection in energy space, therefore to solve it we employed the finite-volume formulation with an upwind scheme of second order. In order to calculate the fluxes at the interface of the cells, we used the piecewise linear method (PLM) with the monotonic central limiter, which is second-order accurate. The intermediate solution was then obtained through the explicit Euler method.

To integrate the spatial diffusion part of the transport equation, we applied the semi-implicit Cranck–Nicolson method, with the gradients calculated at the cell interfaces, using central differences. Therefore, this scheme is second-order accurate.

As a last step, the contributions from the injection term were added, using the Euler explicit method.

The time-steps were chosen in accordance with the Courant–Friedrichs–Lewy (CFL) stability criterion for the minimum time-step of the advection and diffusion equations. Additionally, we imposed the condition that the time-step must be smaller than the time the injector takes to cross one cell.

This paper has been typeset from a T E X/LATEX file prepared by the author.