A Cosmic Zevatron Based on Cyclotron Auto-resonance

Yousef I. Salamin1,2, Meng Wen2,3, and Christoph H. Keitel2,4

1 Department of Physics, American University of Sharjah, POB 26666, Sharjah, UAE; ysalamin@aus.edu
2 Max-Planck-Institute for Nuclear Physics, Saupfercheckweg 1, D-69117 Heidelberg, Germany; wenmeng@hubu.edu.cn, keitel@mpi-hd.mpg.de
3 Department of Physics, Hubei University, Wuhan 430062, People’s Republic of China

Abstract

A Zevatron is an accelerator scheme envisaged to accelerate particles to ZeV energies (1 Zev = 10^{21} eV). Schemes, most notably the internal shock model, have been proposed to explain the acceleration of ultra-high-energy cosmic-ray particles that have been sporadically detected reaching Earth since 1962. Here, the cyclotron auto-resonance acceleration (CARA) mechanism is tailored and used to demonstrate the possible acceleration of particles ejected as a result of violent astrophysical processes such as the merger of a binary system or a supernova explosion. Such events result in emission of highly energetic particles and ultra-intense beamed radiation. In the simultaneous presence of a super-strong magnetic field, the condition for cyclotron auto-resonance may be met. Thus CARA can act like a booster for particles pre-accelerated inside their progenitor by shock waves, possibly among other means. As examples, it is shown that nuclei of hydrogen, helium, and iron-56, may reach ZeV energies by CARA, under which conditions the particles, while gyrating around the lines of an ultra-strong magnetic field, also surf on the waves of a super-intense radiation field. When radiation-reaction is taken into account, it is shown that the ZeV energy gained by a particle can fall off by less than an order of magnitude if the resonance condition is missed by roughly less than 20%.

Unified Astronomy Thesaurus concepts: Gamma-ray bursts (629); Cosmic ray sources (328); Ultra-high-energy cosmic radiation (1733); Interstellar cyclotron emission (835); Compact radiation sources (289); Compact objects (288)

1. Introduction

Cosmic rays are particles which reach Earth from deep space, with energies roughly in the range 10^8–10^{20} eV, and sometimes beyond (Nagano & Watson 2000; Harari 2014). As a result of collisions they make with the atmosphere, showers of secondary particles and flashes of light are produced, which can be seen by detectors on Earth (Linsley 1963; Honda & Honda 2004; Abraham et al. 2007, 2008; Osmanov et al. 2014). The low-energy cosmic-ray particles are believed to originate in active stars and to gain their energies from the shock waves associated with such violent events as supernova explosions (Bell 2013). Particles which arrive with energies toward the end of the above range may be coming from active galactic nuclei with massive black holes at their centers, or from the violent merger of neutron stars. Particles of higher energies are believed to come from extragalactic sources (Allard 2012; Ebisuzaki & Tajima 2014; Fang & Murase 2018).

Extremely rare events of ultra-high-energy cosmic rays (UHECRs) began to be detected more than 50 years ago (Linsley 1963) with energies greater than 10^{18} eV; see (Aloisio 2017) for an excellent review. The flux of cosmic-ray particles gets attenuated as a result of interaction with the cosmic microwave background (CMB) radiation. Cosmic-ray protons, for example, with energy above a minimum of about 4 \times 10^{19} eV, the so-called Greisen–Zatsepin–Kuz'min (GZK) limit (Greisen 1966; Zatsepin & Kuz'min 1966), cannot reach Earth. During their arduous journey through intergalactic space, they lose energy quickly by producing lighter short-lived particles via the Δ^+ resonance, due to interaction with the CMB photons. Examples include

\[ p + \gamma_{\text{CMB}} \rightarrow \Delta^+ \rightarrow p + \pi^0, \quad \text{and} \]

\[ p + \gamma_{\text{CMB}} \rightarrow \Delta^+ \rightarrow n + \pi^+. \]

However, detection of particles heavier than the proton with energies that violate the GZK limit, in its simple form, constitutes no fundamental contradiction (Abbasi et al. 2008, 2014). In fact, measurements by the Pierre Auger Observatory in Argentina (The Pierre Auger Collaboration 2017) suggest that most UHECR particles are nuclei of elements heavier than the proton.

In 1991 the first extreme-energy cosmic-ray (EECR) particle was detected (Abbasi et al. 2008, 2014) with kinetic energy exceeding 3 \times 10^{20} eV. Detection of EECRs meant that either the particles originated in places within the radius dictated by the GZK limit, but have subsequently been accelerated further outside their progenitor by some unknown mechanism, or else that the GZK limit itself had been violated.

The source of EECRs remains a mystery today, despite the existence of models that have been proposed to explain both where the particles come from and what mechanism of acceleration is responsible for the fantastic energies they carry (Hillas 1984; Drury et al. 1994; Chen et al. 2002; Aharonian et al. 2004; Honda & Honda 2004; Drury 2009; Osmanov et al. 2014; Abramowski et al. 2015; Liu et al. 2017; Fang & Murase 2018). This work is concerned with the energy issue, and aims to advance the scheme of cyclotron auto-resonance acceleration (CARA) as a possible resolution to it. Auto-resonance acceleration of electrons and protons, under idealized conditions, has been around for quite some time now (Loeb & Friedland 1986, 1988; Faisal & Salamin 1999; Salamin & Faisal 1998; Salamin et al. 2000, 2015; Galow et al. 2013).
The aim of this paper is to demonstrate that CARA should actually work under known realistic conditions, to be described below, to accelerate cosmic-ray particles to ZeV energies. Realization of this goal may then lead to the development of a more complete astrophysical model in which CARA can be suitably incorporated as a proper mechanism for cosmic-ray acceleration. The paper does not make specific claims about the existence of the non-prohibitive astrophysical conditions where CARA can explain the energy carried by EECRs. It does, however, make the assumption that those particles get pre-accelerated before they enter the region were CARA can give them another tremendous boost.

In CARA, a charged particle gains energy from an ideally coherent, linearly polarized, radiation field monotonically (by multi-photon absorption) if injected along the common direction of its propagation and that of a quasi-uniform magnetic field, and provided an auto-resonance condition is also satisfied to a good approximation. Auto-resonance occurs when the cyclotron frequency of the particle, around the lines of the magnetic field, matches the Doppler-shifted frequency of the radiation field it senses.

The possibility of cosmic-ray acceleration by intense electromagnetic waves has been explored in the past (Gunn & Ostriker 1971), but CARA per se has never been employed to explain acceleration of particles in astrophysical environments, to the best of our knowledge. For it to work for accelerating nuclei to ZeV energies, the scheme requires the simultaneous presence of mega- and giga-tesla magnetic fields, and the highest-power radiation fields believed to exist in the universe. Candidates for such an environment include the polar caps of magnetars, merging neutron stars, and magnetar-powered supernova explosions (Price & Rosswog 2006; Belczynski et al. 2006; Rosswog 2013).

Radiation fields of the required intensity (Aharonian et al. 2008a, 2008b; Hinton & Hofmann 2009; Mészáros 2013; Abramowski et al. 2015; Grenier et al. 2015; Abdalla et al. 2018) may be associated with a gamma-ray burst (GRB). Consider a compact object or a binary-neutron-star merger (Frederiks et al. 2013; Kann et al. 2019) of luminosity $10^{48}$ W. Let half of the output energy (Fan et al. 2006; Abbott et al. 2017) of this event be radiant and beamed (Frail et al. 2001; Granot 2003; Piran 2005) through a circle of radius 100 m, centered on either polar cap. Consequently the emitted radiation in this case can have an intensity around $10^{43}$ W m$^{-2}$.

The general problem of single-particle acceleration, in the simultaneous presence of a radiation field and a uniform magnetic field, will be formulated and its solution revisited in Section 2. The main working equations for CARA will be obtained in Section 3. As examples, acceleration of single protons, and single nuclei of helium and iron, will be investigated in Section 4. The obtained results will be discussed further in Section 5, with emphasis on scenarios that may lead to deviation from the ideal conditions of resonance, radiation loss, and the associated radiation-reaction (RR) effects. Finally, our conclusions will be given in Section 6.

2. Theory

The theoretical background of our investigations will now be outlined, with the aim of making this work self-contained, starting with a specific representation for the electromagnetic fields (Salamin & Faisal 1998; Salamin et al. 2000, 2015; Faisal & Salamin 1999; Galow et al. 2013). Consider a point particle, of mass $M$ and charge $+Q$, injected into a region in which a uniform magnetic field of strength $B_z$ exists parallel to the direction of propagation of a plane-wave linearly polarized radiation field. Employing a Cartesian coordinate system, the combined magnetic and radiation fields may be written, in SI units, as

$$E = \hat{x}E_0 \sin \eta, \quad B = \hat{y}E_0 \frac{\sin \eta + \hat{z}B_0}{c},$$

(3)

(4)

where $\hat{x}$, $\hat{y}$, and $\hat{z}$ are unit vectors in the $x$-, $y$-, and $z$-directions, respectively, $E_0$ and $B_0$ are constants, and $\eta = \omega t - kz$ is the phase of the plane-wave radiation field, of frequency $\omega$ and wavevector $k$ (with $\omega = ck$).

The particle’s relativistic momentum and energy will be denoted by $p = \gamma M c \beta$ and $E = \gamma M c^2$, respectively, where $\beta$ is the velocity of the particle scaled by $c$, the speed of light in vacuum, and $\gamma = (1 - \beta^2)^{-1/2}$ is the Lorentz factor. Neglecting RR at this stage, for simplicity, the relativistic Newton—Lorentz equations (or energy—momentum transfer equations) of the particle, in the above field combination, are

$$\frac{dp}{dt} = Q(E + c\beta \times B),$$

(5)

$$\frac{dE}{dt} = Qc\beta \cdot E.$$  

(6)

Subject to the rather simple initial conditions of position at the origin of coordinates and injection scaled velocity $\beta_0 = \beta_0 \hat{z}$, these equations admit exact solutions, in closed analytic form, for the particle’s trajectory and Lorentz factor. The steps leading to the desired solutions are fairly straightforward. First the $z$-components of Equations (5) and (6) read, respectively

$$\frac{d}{dt}(\gamma \beta_z) = a_0 \omega \beta_z \sin \eta; \quad a_0 \equiv \frac{QE_0}{M \omega c},$$

(7)

$$\frac{d\gamma}{dt} = a_0 \omega \beta_z \sin \eta.$$  

(8)

Note that $a_0$ is a dimensionless radiation field strength, which makes $a_0^2$ a dimensionless intensity parameter. The initial conditions adopted above imply an initial value for the radiation field phase of $\eta_0 = 0$. With this in mind, the left-hand sides of Equations (7) and (8) may be equated and the result integrated to yield a constant of the motion, namely

$$\gamma(1 - \beta_z) = \gamma_0(1 - \beta_0); \quad \gamma_0 = \frac{1}{\sqrt{1 - \beta_0^2}}.$$  

(9)

The analytic solutions may best be arrived at if $\eta$ is employed to replace the time $t$ as a variable (Hartemann et al. 1995) with the following transformation playing a key role:

$$\frac{d}{d\eta} = \omega(1 - \beta_z) \frac{d}{d\eta}.$$  

(10)

In terms of $\eta$ and with the help of (9) the $x$- and $y$-components of Equation (5) now read, respectively

$$\frac{d}{d\eta}(\gamma \beta_x) = a_0 \sin \eta + \alpha(\gamma \beta_x),$$

(11)

$$\frac{d}{d\eta}(\gamma \beta_y) = -\alpha(\gamma \beta_x); \quad \alpha = \frac{QB_0}{M \omega^2 \gamma(1 - \beta_z)}.$$  

(12)
When the constant of the motion expressed by Equation (9) is used, there results

\[ \alpha \to r = \frac{\omega_c}{\omega} \sqrt{\frac{1 + \beta_0}{1 - \beta_0}} ; \quad \text{and} \quad \omega_c = \frac{Q B_s}{M}. \tag{13} \]

Note that \( \omega_c \) is the cyclotron frequency of the particle around the lines of \( B_s \), making \( r \) the ratio of the cyclotron frequency of the particle to the Doppler-shifted frequency of the radiation field sensed by the particle. When Equations (11) and (12) are solved simultaneously, subject to the same initial conditions, they yield

\[ \gamma \beta_x = a_0 \left[ \frac{\cos \eta - \cos (r \eta)}{r^2 - 1} \right], \tag{14} \]

\[ \gamma \beta_z = a_0 \left[ \frac{\sin (r \eta) - r \sin \eta}{r^2 - 1} \right]. \tag{15} \]

These equations give the \( x \)- and \( y \)-components of the particle’s momentum, scaled by \( Mc \). With the help of (9), (10), and (14), Equation (7) may be integrated, subject to same set of initial conditions, to give the \( z \)-component of the scaled momentum

\[ \gamma \beta_z = \gamma_0 \beta_0 + a_0^2 \gamma_0 (1 + \beta_0) \left[ \frac{\sin^2 \eta}{2(r^2 - 1)} \right] + \frac{1 - r \sin \eta \sin (r \eta) - \cos \eta \cos (r \eta)}{(r^2 - 1)^2} \]

\[ \tag{16} \]

Parametric equations will next be derived which give the particle’s coordinates fully analytically. An expression for the \( x \)-coordinate, to begin with, may be obtained with the help of the transformation

\[ \frac{dx}{d\eta} = \frac{\omega \gamma_0 (1 + \beta_0)}{\omega (1 - \beta_0)} = \frac{c}{\omega} \gamma_0 (1 + \beta_0) \gamma \beta_x. \tag{17} \]

Using Equation (14) in Equation (17) and carrying out the integration over \( \eta \), gives an expression for \( x(\eta) \). Expressions for \( y(\eta) \) and \( z(\eta) \) may also be obtained along similar lines. Finally, one gets the following parametric equations for the particle trajectory and Lorentz factor:

\[ x(\eta) = \frac{ca_0}{\omega} \gamma_0 (1 + \beta_0) \left[ \frac{r \sin \eta - \sin (r \eta)}{r^2 - 1} \right], \tag{18} \]

\[ y(\eta) = \frac{ca_0}{\omega} \gamma_0 (1 + \beta_0) \left[ \frac{1 + r^2 \cos \eta - \cos (r \eta) - r^2}{r^2 - 1} \right], \tag{19} \]

\[ z(\eta) = \frac{c}{\omega} \left[ \frac{1 + \beta_0}{1 - \beta_0} \right] \left[ \frac{\beta_0}{1 + \beta_0} \right] \left[ \frac{a_0^2}{3} + \frac{r^2}{4 (r^2 - 1)^2} \right] \eta \\
+ \frac{a_0^2 \sin (2 \eta)}{8 (r^2 - 1)} \\
+ a_0 \left[ \frac{(1 + r^2) \cos (r \eta) \sin \eta - 2r \cos \eta \sin (r \eta)}{(r^2 - 1)^3} \right]. \tag{20} \]

\[ \gamma(\eta) = \gamma_0 \left[ 1 + \frac{a_0^2}{2} (1 + \beta_0) \right. \]

\[ \times \left[ \frac{\cos (r \eta) - \cos \eta)^2 + [r \sin \eta - \sin (r \eta)]^2}{(r^2 - 1)^2} \right]. \tag{21} \]

From Equations (18)–(21) stem all aspects of the particle dynamics in the magnetic and radiation fields expressed by Equations (3) and (4). It ought to be recognized that, viewed as a function of the parameter \( \eta \), the time \( t = \eta / \omega + z(\eta) / c \) is a highly transcendental equation. Nevertheless, Equations (18)–(21) can still be used to investigate evolution in time of the particle’s velocity, momentum, energy, and trajectory during interaction with the given fields. These equations will next be discussed under auto-resonance conditions.

3. CARA: The Working Equations

Equations (18)–(21) have finite limits as \( r \to 1 \). This leads to the much anticipated auto-resonance condition \( r = 1 \). On resonance, the solutions take on the following forms, obtained by taking the limit as \( r \to 1 \) in Equations (18)–(21), respectively:

\[ x(\eta) = \frac{ca_0}{2 \omega} \gamma_0 (1 + \beta_0) [\sin \eta - \eta \cos \eta], \tag{22} \]

\[ y(\eta) = \frac{ca_0}{2 \omega} \gamma_0 (1 + \beta_0) [\eta \sin \eta + 2 \cos \eta - 2], \tag{23} \]

\[ z(\eta) = \frac{c}{\omega} \left[ \frac{1 + \beta_0}{1 - \beta_0} \right] \left[ \frac{\beta_0}{1 + \beta_0} \right] \eta \\
+ \frac{a_0^2}{24} \eta^3 + \frac{a_0^2}{16} \left[ 2 \eta \cos^2 \eta - \sin (2 \eta) \right], \tag{24} \]

\[ \gamma(\eta) = \gamma_0 \left[ 1 + \frac{a_0^2}{8} (1 + \beta_0) \right] \left[ \eta^2 + \sin^2 \eta - \eta \sin (2 \eta) \right]. \tag{25} \]

Equations (22)–(25) contain oscillatory as well as secular terms in the phase variable \( \eta \). The purely secular term in the expression for \( y \) will be responsible for a small transverse drift in the \( y \)-direction. A much more substantial drift in \( z \) is expected due to the first and second terms in Equation (24). More importantly, terms of similar nature in Equation (25) will result in tremendous energy gains by the particle.

It is very important to note that if the resonance condition is met initially, it continues to hold indefinitely. This is guaranteed by the constant of the motion expressed by Equation (9) on account of the fact that \( \beta \approx \beta_z \), for ultrarelativistic particles.

A rough estimate of the kinetic energy of an accelerated nucleus, of atomic number \( Z \) and mass number \( A \), may be obtained from Equation (25). The rapid increase in the kinetic energy, \( K = (\gamma - 1) Mc^2 \), is due mainly to the secular term involving \( \eta^2 \) in Equation (25). For the plane-wave radiation, \( l = c_e \omega_0^2 / 2 \) is the peak intensity, where \( c_e \) is the permittivity of free space. Thus, the particle’s kinetic energy may be calculated from

\[ K(\eta) = K_i + \frac{Q^2}{16 \pi e_0 c_0 M c^2} \left[ I \omega_0^2 \right] \gamma_0 (1 + \beta_0) \gamma \left[ \eta^2 + \sin^2 \eta - \eta \sin (2 \eta) \right], \tag{26} \]
where $K_0 = \frac{(\gamma_0 - 1)MC^2}{2}$ is its initial kinetic energy. Consider interaction for a time equivalent to a change in the radiation field phase from $\eta_0 = 0$ to $\eta = 2n\pi$, where $n$ is the integer (or non-integer) number of phase cycles. In the case $n$ is an integer, only the secular term within the square brackets is nonzero, leading to

$$K(n) = K_0 + \left[ \frac{Q^2}{16\pi^2c_0MC^2} \right] (1-x^2)\eta_0(1 + \beta_0)(2n\pi)^2. \quad (27)$$

On the other hand, averaged over one phase cycle, Equation (26) gives

$$\bar{K} = K_0 + \left[ \frac{Q^2}{16\pi^2c_0MC^2} \right] (1-x^2)\eta_0(1 + \beta_0) \left(1 + \frac{4\pi^2}{3} \right). \quad (28)$$

Approximate versions of Equations (26)–(28) may be obtained by setting $Q = Ze$ and $M = AM_N$, where $M_N \approx (M_n + M_p)/2$ is an average nucleon mass, with $M_n$ and $M_p$ the masses of the neutron and proton, respectively. Thus, the quantity in square brackets becomes $Q^2/(16\pi^2c_0MC^2) \approx 4.071 \times 10^{-28}Z^2/A$, where the expression on the right-hand side is indicated in units of seconds. Note that, because $Z^2/A = 1$ for a proton and a helium nucleus, those two reach roughly the same kinetic energy after interaction with identical numbers of cycles of the same radiation field, but different resonance magnetic fields.

This is quite evident in Figures 1–3. We conclude this section with a note on acceleration efficiency. With the help of Equation (10) and using Equation (9) the rate at which the kinetic energy of the particle changes with time may be written as

$$\frac{dK}{dt} = Ix^2\xi(\eta); \quad \xi(\eta) = \left(\frac{Q^2}{2\pi^2c_0MC^2} \right) \frac{\eta \sin^2 \eta}{\gamma(\eta)} \gamma(\eta). \quad (29)$$

The quantity $\xi(\eta)$ may be thought of as a dimensionless acceleration efficiency parameter, giving the rate at which the radiant energy from the source gets converted to kinetic energy of the particle. This equation may be of use in estimating the energy budgets of potential UHECR sources in order to explain data collected by the Pierre Auger and Telescope Array experiments, assuming the particles have been accelerated by CARA.

### 4. Results

Equations (22)–(25) will now be used in accurate one-resonance single-particle calculations. Meeting the resonance condition can be a delicate matter. So, part of this section will be devoted to investigating dependence of the end results, of interest to us in this work, on fluctuations around resonance. The radiation loss and RR effects will also be discussed.

#### 4.1. Examples

Single-particle dynamics, in several specific magnetic and radiation field environments, will be investigated in this section, on the basis of Equations (22)–(25). Typically, a particle gyrates around the lines of the magnetic field, and follows a semi-helical trajectory of increasing cross-sectional area, as Figure 1(a) aims to demonstrate. For the parameters used in this example, the maximum transverse extension of the semi-helical trajectory is about 0.2 m, while the total axial excursion is over 3200 km. So, the particle’s path is essentially a straight line. On the other hand, the particle’s kinetic energy increases monotonically, by continuous multi-photon absorption from the radiation field. For the parameters employed, a proton’s kinetic energy reaches 1.5 ZeV, as shown in Figures 1(b) and (c). More specifically, the energy reaches PeV, EeV, and ZeV in distances of several millimeters, tens of meters and thousands of kilometers, respectively. The trajectory and evolution of the kinetic energy are shown for interaction with six phase $\eta$-cycles (one phase cycle $= 2\pi$). The kinetic energy curve exhibits 12 *knee*s, each of which representing a *kick* in the particle’s energy, due to interaction with one-half of a phase cycle. Figure 1(a) also shows a semi-helical path of six windings. Each winding is the result of interaction with one complete phase cycle. There exists a one-to-one correspondence between the kicks in Figure 1(b) and the windings of the semi-helical trajectory (each winding corresponds to two successive kicks).

Figure 1 is meant to illustrate CARA. The magnetic field strength needed to achieve auto-resonance, in this particular example, is $B_x = 1.92753 \times 10^1$ T, the likes of which are believed to be associated with classical pulsars (Reisenegger 2001). In Figures 2 and 3, the parameter values employed could be associated with magnetar-powered supernovae, according to recent studies (Greiner et al. 2015; Sukhbold & Woosley 2016; Sukhbold Salamin, Wen, & Keitel...
Next, the seemingly sensitive dependence of the kinetic energy of a specific particle species, identified by its charge-to-mass ratio, \( Q/M \), on variations, \( \Delta r \), around the resonance value of \( r = 1 \), will be investigated. Since little has been particularly specific so far in this work about the magnetic and radiation field environments in which the particles are accelerated, only general statements can be made about the parameter space centered about the values which lead to resonance. As can readily be seen in Equation (13), the resonance condition depends on the static magnetic field \( B_s \), the injection speed \( \beta_0 \) (or, equivalently, the scaled injection energy \( \eta \)), and the radiation frequency \( \omega \) (or, equivalently, the wavelength \( \lambda \)). Holding \( Q/M \) fixed, Equation (13) yields

\[
\frac{\Delta r}{r} = \frac{\Delta B_s}{B_s} + \frac{\Delta \omega}{\omega} + \gamma_0^2 \Delta \beta_0,
\]

where the symbol \( \Delta X \) stands for a spread in the possible values of the parameter \( X \in \{B_s, \omega, \beta_0\} \) around its on-resonance value. Thus, variations in any one, or more, of the parameters in this set lead to a departure from \( r = 1 \). This means that investigation of the dynamics over a set of values of \( r \) around unity is tantamount to studying the dynamics over many sets of different values of these quantities simultaneously. As a first example, variations in the magnetic field only will be considered. Keeping \( \beta_0 \) and \( \omega \), fixed, values of \( r \) are varied for the three nuclear species \( H^+, He^+, \) and \( Fe^{26+} \). If one lets \( r \rightarrow 1 + \Delta r \) in Equations (18)-(21) the particle dynamics may be investigated over values of \( r \in \{1 - \Delta r, 1 + \Delta r\} \), for the given charge-to-mass ratios. Viewed as a function of \( \Delta r \), the Lorentz factor, and hence the exit kinetic energy, has oscillatory (sinusoidal) terms. That should explain the bumps on both sides of the main peak in Figure 5 below.

In Figure 4, variations in the particle’s exit kinetic energy are shown as a function of the detuning \( \Delta r \) around resonance (\( r = 1 \)) with the interaction time (expressed indirectly through \( \eta \)) and also with the distance over which the particle interaction with the magnetic and radiation fields takes place. Figure 5 shows the kinetic energy of each particle, at the end of its interaction with the given on-resonance magnetic field and four phase cycles, as a function of the detuning \( \Delta r \).

Note that Equation (13) implies that \( \Delta r/r = \Delta B_s/B_s \), for fixed values of \( \omega \) and \( \beta_0 \). Using the magnetic field values that correspond to resonance, shown in Figure 3, leads to the conclusion that a 40% deviation, \( \Delta B_s \), from its corresponding resonance values, leads to a similar percentage deviation from resonance (detuning \( \Delta r \)). According to Equation (30) similar conclusions may be arrived at when variations, \( \Delta \omega \), around the resonance value of \( \omega \), are considered, keeping \( B_s \) and \( \beta_0 \) fixed.

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4 We note that the resonance condition may be violated by other effects such as particle–particle collisions, as well as lack of coherence and other deviations from our assumptions about the radiation field. While those issues need to be investigated in future publications, our investigations in the present work, based on the parameter set \( \{B_s, \omega, \beta_0\} \), already hint clearly that small deviations from resonance are not very critical for explaining cosmic ZeV acceleration by CARA.
at their respective values on resonance. Figures 5(c) and (d) indicate that only the parts of the spectrum around resonance contribute to the acceleration of ZeV particles in broad frequency radiation fields like a thermal spectrum. Finally, keeping $B_s$ and $\omega$ fixed at their respective on-resonance values and varying those of $\beta_0$ in the interval $[-0.4/\gamma_0, 0.4/\gamma_0]$ results in $\Delta r \sim 40\%$, according to Figure 5. Numerical calculations based on Equation (21) indicate that the kinetic energy at the end of the interaction, in all cases considered, remains roughly within one order of magnitude of its on-resonance maximum, as the values of $r$ are detuned by about 10%–20% on either side of $r = 1$. Off-resonance conditions cause the particle to phase-slip behind the radiation wave and gain little energy from it, if at all.

4.3. Radiation Loss

It was remarked, introducing Figure 1(a), that the trajectory of a particle accelerated by the CARA scheme can essentially be considered linear. It is also well known that, in linear accelerators on Earth, radiation loss by the accelerated particle is negligibly small in general. The instantaneous power radiated by the particle during acceleration is given by the relativistic generalization of the Larmor formula (Jackson 1998)

$$ P = \frac{Q^2 \gamma^6}{6\pi e_0 c} \left[ \left( \frac{d\beta}{dt} \right)^2 - \left( \beta \times \frac{d\beta}{dt} \right)^2 \right]. \quad (31) $$

The presence of the factor $\gamma^6$ upfront in this expression can make the radiated power quite large, even exceeding the rate of energy gain, $d\mathcal{E}/dt$, beyond some point. To avoid running into numerical instabilities, the following alternative expression for the power, in terms of derivatives with respect to the phase $\eta$ and the other on-resonance parameters, will be used (Salamin et al. 2000):

$$ P(\eta) = \frac{Q^2 \gamma^2}{6\pi e_0 c} \left[ \left( \frac{d\beta}{d\eta} \right)^2 + \gamma^2 \left( \beta \cdot \frac{d\beta}{d\eta} \right)^2 \right]. \quad (32) $$

Power lost through radiation will be assessed, in relation to the rate at which energy is gained by the particle, by investigating the dimensionless ratio

$$ R(\eta) = \frac{1}{P} \frac{d\mathcal{E}}{dt}. \quad (33) $$

This ratio is shown as a function of the phase variable $\eta$ in Figure 6, for the cases considered in Figures 3–5. According to Figure 6, $R \gg 1$ during interaction with a small fraction of the first phase cycle. The corresponding interaction time over which the spike in $R$ happens is less than 10 ps, according to the inset in Figure 8, to be introduced below. This happens while the initial part of the particle’s trajectory is still almost straight, implying a relatively small radiation energy loss rate. Subsequent to this brief spike in the values taken up by this ratio, $R$ drops sharply and may approach unity (or less) only over brief intervals during the particle’s motion, as may be inferred from the insets in Figure 6.
4.4. Radiation-Reaction Effects

At the time of emission of the radiation, the particle recoils as it experiences a RR force, which must be taken into account as, especially in circumstances for which \( R \to 1 \) (Hadad et al. 2010; Tamburini et al. 2010, 2011; Seto et al. 2012; Poder et al. 2018; Tursunov et al. 2018). To assess the RR effects quantitatively, Equations (5) and (6) will now be replaced by (Hadad et al. 2010; Seto et al. 2012)

\[
\frac{dp}{dt} = Q(E + c\beta \times B) - f_1 + f_2 + f_3, \tag{34}
\]

\[
\frac{dE}{dt} = Qc\beta \cdot E - P_1 + P_2 + P_3, \tag{35}
\]

respectively. In these equations, \( f_1 \) is a force term and \( P_i \) is a power term. With \( i = 1, 2, \) and 3, the force and power terms are given by

\[
f_1 = \frac{\beta}{c} P_i, \tag{36}
\]

\[
f_2 = \gamma A \left[ \frac{Q}{Mc^2} \right] \left[ \frac{\partial}{\partial t} + c\beta \cdot \nabla \right] (E + c\beta \times B), \tag{37}
\]

\[
f_3 = A \left[ \frac{Q}{Mc^2} \right] \left[ (\beta \cdot E) \frac{E}{c} + (E + c\beta \times B) \times B \right], \tag{38}
\]

\[
P_1 = \gamma^2 A \left[ \frac{d\beta}{d\eta} \right]^2 + \gamma \left( \beta \cdot \frac{d\beta}{d\eta} \right)^2 \omega_c^2 = P(\eta), \tag{39}
\]

\[
P_2 = \gamma A \left[ \frac{Q}{Mc^2} \right] \left[ \frac{\partial}{\partial t} + c\beta \cdot \nabla \right] (c\beta \cdot E), \tag{40}
\]

\[
P_3 = A \left[ \frac{Q}{Mc^2} \right] (E + c\beta \times B) \cdot E; \quad A = \frac{Q^2}{b\pi c^2}. \tag{41}
\]

This classical treatment of RR is assumed to be a good approximation, in light of the fact that \( \chi \), the quantum efficiency parameter, can be quite small (Hadad et al. 2010; Tamburini et al. 2010). This parameter is defined by (Ritus 1985):

\[
\chi = \frac{Q\hbar \gamma}{M^2c^3} \sqrt{(E + c\beta \times B)^2 - (\beta \cdot E)^2}, \tag{42}
\]

where \( \hbar \) is the reduced Planck constant. For an electron of mass \( M = m_e \), the RR effects can be quite sizeable. But for protons and heavier nuclei (mass \( M \gtrsim 2000m_n \)) of the same \( \gamma \)-factor, \( \chi \) is at least six orders of magnitude smaller. For all of the examples considered, and for the parameters used, \( \chi \ll 1 \), as Figure 7 clearly demonstrates.

Note that \( f_2 \) and \( P_2 \) are a factor of \( \gamma \sim 10^{12} \) times smaller than \( f_1 \) and \( P_1 \), respectively, for \( Z \) \& \( N \) particles. Furthermore, \( f_3 \) and \( P_3 \) are a factor of \( \gamma^2 \sim 10^{24} \) times smaller than \( f_1 \) and \( P_1 \), respectively. Consequently, only \( f_1 \) and \( P_1 \) will be retained in Equations (34) and (35) which will subsequently be solved numerically. For ultrarelativistic particle energies and small field gradients, the dropped terms contribute negligibly. Results for the exit kinetic energies, stemming from those solutions, are displayed with dotted lines in Figure 2. The effect of RR on the radiated power and energy gain by the nuclei of \( H^1 \), \( He^4 \) and \( Fe^{56} \) is to lower those quantities by roughly less than one order of magnitude. Another notable related effect is shortening of the total excursion distance, in most cases considered, due to RR. This is reminiscent of the effect of friction on the motion of macroscopic objects.

Similar plots are displayed in Figure 3. With the RR effects taken into account, all three particles are shown to leave the loosely defined interaction region with kinetic energies in excess of 10^{20} \text{eV}. According to Figure 3, the RR effects are responsible for a reduction in the exit kinetic energy. A reduction in the particle’s excursion distance, during the same interaction cycle and with RR taken into account, is also clear in Figures 3(a) and (b).

Up to this point, time dependence has been expressed indirectly in terms of the dependence upon the phase \( \eta \). For the examples of Figure 3, variation of the time \( t \) with \( \eta \) is shown in Figure 8, according to which the interaction times of the nuclei of \( H^1 \), \( He^4 \) and \( Fe^{56} \) are, roughly, 96 \text{ms}, 24 \text{ms}, and 20 \text{ms}, respectively. In terms of the period \( \tau = \lambda/c \) of the GRB...
employed, these are equivalent to $5.76 \times 10^{14} \tau$, $1.44 \times 10^{14} \tau$, and $1.2 \times 10^{14} \tau$, respectively.

It cannot be concluded with certainty, based on Figures 2 and 3, that the RR effects act to universally lower the accelerated particle kinetic energy so drastically in CARA. With RR taken into account according to the approximate version of Equations (34) and (35) the right-hand sides of (7) and (8) are no longer the same. Subsequently, subtraction of the RR-based equivalent to (7) from the approximate RR-equivalent to (8) does not lead to a constant of the motion analogous to that expressed by Equation (9) but to

$$\frac{d}{dt} \gamma(1-\beta) = \frac{(1-\beta)P(t)}{Mc^2},$$

instead. Making the $t \rightarrow \eta$ transformation in Equation (43) and integrating the result formally over $\eta$ leads to

$$\gamma(1-\beta) = \gamma_0(1-\beta_0) - \frac{J}{\omega Mc^2},$$

in which

$$J(\eta) = \int_0^\eta P(\eta')d\eta'.$$

Using Equation (44) in the second of Equation (12), with the replacement $B_\delta = B_{\delta\text{res}} + \Delta B_\delta$, where $B_{\delta\text{res}}$ is the on-resonance value of $B_\delta$ when RR is neglected, results in

$$\alpha = \frac{Q(B_{\delta\text{res}} + \Delta B_\delta)}{M \omega [\gamma_0(1-\beta_0) - J/(\omega Mc^2)]} < \frac{Q(B_{\delta\text{res}} + \Delta B_\delta)}{M \omega \gamma_0(1-\beta_0)} = r + \frac{Q \Delta B_\delta}{M \omega D},$$

where $\omega_D$ is the Doppler-shifted frequency of the radiation wave.

In the analytic discussions of CARA in the absence of RR, $r$ plays the role of a resonance parameter. By analogy, that role would most probably be played, albeit roughly, by the quantity $\alpha$ defined by Equation (46) when RR is taken into account. Ideally, this should emerge from analytic solutions to Equations (34) and (35). In the absence of such solutions, $\alpha$ may only be used to hint at the existence of a resonance condition different from $r = 1$.

Having been derived from the resonance condition in the absence of RR, Equation (30) can no longer be relied upon in the presence of RR. In the subsequent analysis of the acceleration process with RR, total dependence upon the numerical solutions outlined above is inevitable.

To further assess departure from the $r = 1$ resonance due to inclusion of RR, Figure 9 has been produced by scanning the static magnetic field values over $\Delta B_\delta/B_\delta \in [-0.4, 0.4]$ centered around the on-resonance values employed in Figure 3. This figure may be read in conjunction with Figure 4, but with care. In Figure 4, $\Delta B_\delta/B_\delta = \Delta r/r$, which does not hold when RR is taken into account. In the absence of RR, maximum ZeV energy is reached for parameters which make the detuning $\Delta r = 0$, as in Figure 4. However, when RR is included in the numerical calculations, the maximum ZeV energies are shown in Figure 9 to be reached for magnetic fields weaker than the on-resonance values when RR is ignored.

In all cases considered, the reachable maximum ZeV energies in the presence of RR are less than in its absence. Effects of RR on the exit kinetic energy are summarized in Table 1. The entries are based on the results displayed in Figures 4 and 9. In Table 1, the deviation is defined by $\Delta K = K - K'$, where $K(K')$ is the kinetic energy without (with) RR. Further support for these conclusions is presented in Figure 10. For every parameter $X \in \{B_\delta, \omega, \beta_0\}$, Figure 10 shows clearly that the reachable maximum energy is smaller than when RR is neglected, and that the maximum energy is reached, in each case, for a set of RR-based parameters different from those related by $r = 1$, without RR. The maximum attainable energies of the nuclei of helium are lowered by RR less than both hydrogen and iron nuclei, in
agreement with the results shown in Figure 3. More importantly, the optimal exit kinetic energies can be smaller by much less than an order of magnitude, when the RRs are taken into account, than when resonance is met without RRs.

Besides the impact, shown in Figure 10, on them when the RR effects are taken into consideration, the exit kinetic energies are lowered due to deviation from the resonance values of ω and β0. This observation leads to conclusions that agree qualitatively with those based on a spread, ΔB∗, about the corresponding on-resonance value of B∗. Results without RR (solid lines) and with RR (dotted lines) are shown together in each respective case, to allow for comparisons to be made. ZeV particles are obtained for values of the parameters X′ ∈ [B′, ω′, β′] in windows ΔX′ ∼ 10%. For example, scanning the static magnetic field beyond the range between B′1 = 0.1188B∗ and B′2 + 0.311B∗ leads to proton energies up to one order of magnitude lower than the maximum value obtained for B′ = 0.844B∗. When RR is taken into account, the maximum kinetic energies are obtained for weaker static magnetic fields, as well as at smaller values of ω and β0. For more accurate comparisons, see Table 2, in which differences are shown between the on-resonance values of X (without RR) and the values X′ that correspond to maximum energy gain when RR is taken into account. For example, the energy for a proton peaks at a value of B′ that is about 11.8% below the corresponding on-resonance value of B∗.

Thus, by including the RR effects, the realistic astrophysical conditions have been better simulated in our calculations. It has been demonstrated that, under these conditions, particles in the ZeV range can be generated, even for parameters that violate the resonance condition substantially.

Depletion may occur during acceleration inside the source, as well as en route to Earth. During acceleration, a particle typically interacts with a photon of Doppler-shifted frequency f′0(1 − β0), where the constant of the motion Equation (9) has been used, and f = c/λ. The photon energy in the rest frame of the particle is then ϵph = (hc/λ)γ0(1 − β0). For the parameters of Figures 2 and 3, one calculates ϵph ∼ 0.285 eV and ∼24.8 eV, respectively. These energies lie well below the threshold of every one of the following processes (which, otherwise, would contribute to depletion): (a) e+e− pair-production at ϵph ≥ 1.022 MeV, (b) nuclear excitation and disintegration at ϵph ≥ 8 MeV (the average nuclear binding energy), and photodisintegration above the photomeson threshold ϵph ≥ 140 MeV, and so on (Biehl et al. 2018; Morejon et al. 2019). Other photodisintegration requires much more energetic photons. For example, photodisintegration becomes relevant for γ0 = 500 (Figure 3) at a GRB wavelength λ ≈ 1.55 × 10−16 m (f ≥ 1.93 × 1019 Hz).

During propagation, the particles interact with (a) the CMB radiation, (b) the extragalactic background light, and (c) the intergalactic magnetic fields (Aloisio 2017). The GZK suppression of protons occurs via photodisintegration at particle energies ≥5 × 1019 eV. Also, proton suppression by pair-production begins to kick in below a threshold of ~2 × 1018 eV (Aloisio et al. 2013a, 2013b). Most recent observations favor heavy nuclei in the UHECR spectrum (Aab et al. 2020a, 2020b). A drop in the mass spectrum of a nucleus of mass number A occurs at ~A × 1018 eV, due to photodisintegration on far-infrared photons, and to pair-production (Stecker et al. 2006; Allard 2012; Kneiske et al. 2004). A second drop occurs near ~A × 4 × 1018 eV, due to photodisintegration on CMB photons. However, these studies conclude that depletion of UHECR nuclei by photodisintegration becomes relevant for energies of ~1021 eV, even for light nuclei, and can then be safely omitted in propagation calculations, according to Allard (2012).

5. Discussion

The process of single-particle energy gain in CARA will be highlighted further in this section. This will be followed by a discussion of the astrophysical conditions that may lead to ZeV gain.

5.1. The Underlying Process of Energy Gain

A particle gains zero energy from interaction with an integer number of cycles of a plane-wave radiation field, in the absence of boundaries or a material medium, and under off-resonance conditions, according to the Lawson–Woodward theorem (Lawson 1979). Effectively, the particle gains energy during interaction with the positive (accelerating) half of a single cycle, only to lose it entirely during interaction with the following (identical) negative (decelerating) half. In the added presence of a static magnetic field B∗, however, the Lawson–Woodward theorem gets circumvented and a resonance condition may be met. On resonance, the electric field vector of, say, a circularly polarized plane wave, gyrates about the direction of propagation at the angular frequency, ω, of the wave. The particle, while in cyclotron motion about the lines of B∗, senses the radiation field at the Doppler-shifted frequency ωD = ω√(1 − β0)/(1 + β0). The argument holds just as well for a linearly polarized radiation field, which may be represented in terms of a superposition of two circularly polarized ones, of opposite helicity. In either case, autoresonance guarantees that the vectors β and E in Equation (8) maintain the same orientation relative to each other. This means the rate dEdt will remain positive, implying energy gain, during interaction with both halves of every radiation field phase-cycle. Thus, energy of the particle continues to increase monotonically. Barring any unaccounted for perturbations, the resonance condition continues to hold true, due to the constant of the motion expressed by Equation (9) and the particle continues to be accelerated.

5 Our discussion has focused on the linear regime of single-photon absorption, due to the energy gain over half a cycle of the radiation field with strongly Doppler-extended wavelength (and period in time). In the multiphoton regime many photons can be absorbed simultaneously and the situation could be different. The photodisintegration effects may be negligible in our examples, even in the presence of additional, moderately intense, background photon fields.

6 In order to stay focused on CARA as a new acceleration mechanism, the statements above have been made despite the lack of knowledge of the various contributing/competing pre-acceleration scenarios for which CARA would serve as a booster. One should be aware that lack of knowledge of the involved distances and complex environmental background still leave room for considerable uncertainty.
5.2. Astrophysical Environment for Acceleration

Chances that the parameters $\beta_0, \omega, B_s, \text{ and } Q/M$, will conspire to satisfy the auto-resonance condition can surely be very small. Nevertheless, this should not be terribly discouraging, in light of the fact that EECR events are quite rare, indeed. For example, the Telescope Array experiment (Abbasi et al. 2014) reported detecting 72 events only, with energies of $57 \ EeV$ or more ($1 \ EeV = 10^{18} \ eV$) over the five-year period between 2008 and 2013. It has been estimated that EECRs may be detected on Earth at a rate of one particle per kilometer squared per century (Nagano & Watson 2000).

This work has not been specific about any known astrophysical environment where conditions for CARA can initially be met. One possibility can be a small region centered on either pole of a compact object, where the right magnetic and radiation fields may be found. To show that the numbers employed in the examples discussed above are not totally

![Figure 10](image-url)

**Figure 10.** RR effects shift the conditions of reaching optimal energies, from resonance expressed by $r = 1$, which holds when the RR effects are neglected. Shown here is variation of the end-of-interaction particle kinetic energy which results, in each case, from a deviation, $\Delta X$, from the corresponding resonance value, determined by $r = 1$, of the parameter $X \in \{B_s, \beta_0, \omega\}$. The central values, like in Figure 3, are: $\omega = \frac{2\pi c}{\lambda}$ with $\lambda = 5 \times 10^{-11} \ m$, $\gamma_0^2 \beta_0 = \gamma_0 \sqrt{\gamma_0^2 - 1}$, with $\gamma_0 = 500$, and $B_s = 393.3 \ \text{MT, 781.2 \ \text{MT, and 840 \ \text{MT}}$, for nuclei of $H^\ +, \text{He}^\ +_2$, and $\text{Fe}^\ +_26$, respectively. Solid lines: without RR (as in Figures 4 and 5); dotted lines: with RR (based on Figure 9).

| Off-resonance Parameter Deviations for Moderate Energy Drop | $H^\ +$ | $\text{He}^\ +_2$ | $\text{Fe}^\ +_26$ |
|------------------------------------------------------------|--------|---------|--------|
| $B_s^/B_s \ (%)$                                           | 84.8   | 94.4    | 74.8   |
| $\Delta B_s/B_s \ (%)$                                    | [-11.8, 31.1] | [-16.1, 21.2] | [-8.6, 51.9] |
| $\omega/\omega \ (%)$                                     | 82.0   | 92.4    | 71.2   |
| $\Delta \omega/\omega \ (%)$                              | [-11.7, 29.3] | [-15.6, 22.5] | [-15.7, 46.6] |
| $\gamma_0^2 \beta_0^2 \ (%)$                              | 83.6   | 95.2    | 71.6   |
| $\gamma_0^2 \Delta \beta_0^2 /\beta_0^2 \ (%)$            | [-12.4, 58.9] | [-16.8, 21.4] | [-16.2, 82.7] |

**Note.** Values of the parameters $X' \in \{B_s', \omega', \beta_0'\}$ for which optimal exit kinetic energies are obtained in the presence of RR (read from the data displayed in Figure 10) compared to the corresponding on-resonance values of the parameters $X \in \{B_s, \omega, \beta_0\}$, calculated from the condition $r = 1$ (without RR). Shown here also are detuning windows $\Delta X'$ around the values $X'$ in which particle energies are reduced by at most one order of magnitude as compared to the maximum value when the RRs are included.

5.2. Astrophysical Environment for Acceleration

Chances that the parameters $\beta_0, \omega, B_s$, and $Q/M$, will conspire to satisfy the auto-resonance condition can surely be very small. Nevertheless, this should not be terribly discouraging, in light of the fact that EECR events are quite rare, indeed. For example, the Telescope Array experiment (Abbasi et al. 2014) reported detecting 72 events only, with energies of $57 \ EeV$ or more ($1 \ EeV = 10^{18} \ eV$) over the five-year period between 2008 and 2013. It has been estimated that EECRs may be detected on Earth at a rate of one particle per kilometer squared per century (Nagano & Watson 2000).

This work has not been specific about any known astrophysical environment where conditions for CARA can initially be met. One possibility can be a small region centered on either pole of a compact object, where the right magnetic and radiation fields may be found. To show that the numbers employed in the examples discussed above are not totally
improbable, consider for instance that highly powerful GRBs (Kumar & Zhang 2015; Wang et al. 2018; Guessoum et al. 2018) can have isotropic energies of the order of $10^{46}$–$10^{48}$ erg, in addition to being accompanied by the ejection of protons and other particles, pre-accelerated by internal shock waves to relativistic speeds. The radiant energy is mostly beamed into a jet of a few-degree half-angle, and is typically given off over a period of time that can be as small as a fraction of a second to a few seconds (and as large as 17 minutes in rare cases). Thus the radiation-field intensities of $10^{38}$–$10^{44}$ W m$^{-2}$, used in our examples, may be considered reasonable.

The magnetic field strength in such an environment, and its direction at, and close to, the pole, can be right for the CARA scheme. However, away from the polar caps of a compact object, the magnetic field lines can be extremely curved and the radiation field intensities fall way below what is needed for CARA to work. Typically, $B_{z} \sim 1/\zeta_{z}$, for $z >$ radius $\sim$10 km (Belczynski et al. 2006; Price & Rosswog 2006; Shapiro & Teukolsky 2007; Rosswog 2013).

During the relatively brief interaction of the particle with the magnetic and radiation fields, Equation (6) implies that, if met initially, the resonance condition continues to hold so long as the conditions outlined in Section 2 are met. Radiation in the astrophysical environments of relevance to this work can be coherent (Huege & Falcke 2003; Gainullin & Zlobin 2005; Ioka 2005). On the other hand, a recent study has shown that particles can be efficiently accelerated by incoherent radiation in the wakefield of a plasma (Benedetti et al. 2014).

6. Conclusions

Basic elements of the scheme of cyclotron auto-resonance acceleration (CARA) have been reviewed in this paper and tailored to meet astrophysical conditions of ultra-strong magnetic and super-intense radiation fields, which would work to efficiently accelerate protons and heavier nuclei to ZeV energies. The results include equations giving the trajectory and Lorentz factor of a single particle, in closed analytic form. This has been accomplished in the absence of RRs. RRs have later been approximately included via numerical simulations and their effects on the optimal energy gains have been assessed.

Mega- and giga-telsa magnetic fields, and radiation fields of intensity $10^{38}$–$10^{44}$ W m$^{-2}$, have been shown to accelerate nuclei of hydrogen, helium, and iron to ZeV energies, over distances ranging from several hundred meters to many kilometers. Directions of the magnetic field and that of propagation of the radiation have been assumed to be strictly parallel. Magnetic field lines through a small region around the pole (within the polar cap) of a compact object, for example, may be straight and strong enough over a long distance, but can be severely curved elsewhere.

The particles have also been assumed to be pre-accelerated inside their source to relativistic speeds before entering the region for further acceleration by CARA.

Key for the process to work is meeting the auto-resonance condition. It has been shown that, when RRs are neglected, once the resonance condition is met initially, it would hold almost exactly throughout. In this regime, dependence of resonance upon the initial conditions and magnetic and radiation field parameters has been investigated. It has been shown that the (non-resonance) ZeV energy gained by a particle stays within an order of magnitude of that maximum value if resonance is slightly missed due to a spread of a few tens of percentage points in the values of one or more of the parameters $B_{z}$, $\omega$, or $\beta_{0}$. It has been demonstrated via numerical calculations that the ZeV energy gains are still maintained, also to within an order of magnitude, when RRs are taken into account.

CARA has been advanced as a possible mechanism for UHECR acceleration without reference to an astrophysical environment known with any certainty to host the needed backgrounds for resonance to occur. The assumption has been made throughout that the particles are pre-accelerated inside the source before they are injected for further acceleration by CARA. Pre-acceleration inside the progenitor may be due to any or all of the well investigated mechanisms of Fermi acceleration at shock waves, magnetic reconnection, or unipolar induction. As such, CARA complements those models and works as a booster mechanism.

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Orcid iDs

Yousef I. Salamin https://orcid.org/0000-0003-2343-4031
Meng Wen https://orcid.org/0000-0001-7567-5350
Christoph H. Keitel https://orcid.org/0000-0002-1984-1470

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