Heavy nuclei at the end of the cosmic ray spectrum?

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We provide an account of the possible acceleration of iron nuclei up to energies ~ 300 EeV in the nearby, metally-rich starburst galaxy NGC 253. It is suggested that particles can escape from the nuclear region with energies of $\sim 10^{15}$ eV and then could be reaccelerated at the terminal shock of the galactic superwind generated by the starburst, avoiding in this way the photodisintegration expected if the nuclei were accelerated in the central region of high photon density. We have also made estimates of the expected arrival spectrum, which displays a strong dependency with the energy cutoff at the source.

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I. INTRODUCTION

The discovery of extensive air shower events with energies above 100 EeV (see Ref. [1] for a recent survey) confirms that the cosmic ray (CR) spectrum does not end with the expected Greisen-Zatsepin-Kuz'min (GZK) cutoff [2]. The origin of this GZK cutoff is energy degradation of the CR particles (usually assumed to be nucleons and nuclei) by resonant scattering processes with the diffuse background radiation that permeates the universe. The observed tail of the spectrum could be, consequently, originated in a bunch of nearby sources [3].

Preferred sites for proton acceleration are astrophysical scenarios where large-scale shocks occur, as for instance in the hot spots of powerful radio-galaxies [4]. It could appear that, because of the high-energy cutoff of shock acceleration increases with the charge number of the nucleus, heavy ions would be nice candidates for ultra high-energy CRs. However, AGNs and radio galaxies are widely thought as regions of very low metallicity and, in addition, it is well established that above 200 EeV nuclei should be photodissociated by the 2.73 K photon background in a few Mpc [5]. Thereupon, nuclei acceleration up to the highest energies within astrophysical environments is seldom considered in the literature.

Nonetheless, there has been a recently renewed interest in the propagation of heavy nuclei [5–9]. This renewal is mainly sustained by two facts: i) a medium mass nucleus is the particle that provides the best fits of the shower development of the highest energy CR event [10], and ii) the arrival direction of such event roughly points towards the nearby metally-rich galaxy M82 [11].

Despite the aforementioned studies on nuclei propagation, it is far from clear whether iron or other heavy nuclei can be accelerated up to energies ~ 300 EeV in starburst galaxies like M82. One could naively expect that, since the size scale of the starburst region is of the order of the gyroradius of 300 EeV–($Z = 26$) nuclei, strong shocks could diffusively accelerate these ions to ultra-high energies. But this hope fades away as soon as one notices the large photon energy densities (mostly in the far infrared) measured in the central regions of these kind of galaxies: iron nuclei are photodisintegrated long before they can reach the required Lorentz factors.

We shall argue in this paper that, despite the mentioned problem, iron nuclei can be actually acclerated in nearby starburst galaxies up to energies ~ 300 EeV if a two-step process is involved. The crucial point is that for energies above $\sim 10^{15}$ eV, acceleration occurrs in the terminal shock of the starburst superwind, well outside the problematic nuclear region. Since NGC 253 is a southern object which has been scarcely discussed in relation to CRs (for a brief discussion of M82 as CR accelerator see [11]), we shall focus on it for our quantitative estimates. However, we emphasize that due to the similarity between both galaxies, our conclusions will be correct to M82 within the order of magnitude. Let us start with a recapitulation of some observational features of NGC 253.

II. THE STARBURST GALAXY NGC 253

Starbursts are galaxies undergoing a massive and large-scale star formation episode. Their characteristic signatures are strong infrared emission (originated in the high levels of interstellar extinction), a very strong HII-region-type emission-line spectrum (due to a large number of O and B-type stars), and a considerable radio emission produced by recent supernova remnants. Typically, the starburst region is confined to the central few hundreds of parsecs of the galaxy, a region that can be easily 10 or more times brighter than the center of normal spiral galaxies.

NGC 253 has been described as the archetypal starburst galaxy by Rieke *et al.* [12], and as a prototype of superwind galaxy by Heckman and collaborators [13]. This object, whose distance is estimated in the range 2.5 - 3.4 Mpc [14,15], has been extensively studied from radio to γ-ray wavelengths [16–18]. More than 60 idividual compact radio sources have been detected within the central 200 pc of the nuclear region of NGC 253 [19], most of which are supernova remnants (SNRs) of only a few hundred years old. According to estimates from observations at different frequencies the supernova rate is as high as $0.2 - 0.3$ yr^{-1} [19,20].

The central $\sim$ 80 pc of the galaxy contain around 24.000 O and 3.000 red supergiant stars [20], in addition to the SNRs and numerous HII regions. This means that the massive star formation rate is ~ 0.1 M$_\odot$ yr^{-1}. Strong [Fe II] emission has been also detected with a total [Fe II] (1.644 μm) luminosity of $\sim 2.8 \times 10^{39}$ erg s^{-1}, which reflects the very rich iron production in the supernovae and their associated shocks [21].

In the light of such a concentrated activity it is not surprising that strong physical, morphological, and kinematic evidence for the existence of a galactic superwind has been found in NGC 253 [22,13]. Galactic-scale superwinds are driven by the collective effect of supernovae and massive star winds. The high supernovae rate creates a cavity of hot gas ($\sim 10^8$ K) whose cooling time is much greater than the expansion time scale. Since the wind is sufficiently powerful, it can blow out the interstellar medium of the galaxy avoiding to remain trapped as a hot bubble. As the cavity expands a strong shock front is formed on the contact surface with the cool interestellar medium. Shock interactions with low and high density clouds produce X-ray continuum and optical line emission, respectively, that has been directly observed [22]. In addition, kiloparsec regions well outside the disk present double emission-line profiles with line splitting of 200-600 km s^{-1}, a clear evidence of mass motion. The morphology of the optical emission line nebulae indicates that the outflowing gas is located along the walls of a cone that is limb-brightened, typical of a superwind in a blowout phase.

The shock velocity can reach several thousands of kilometers per second and ions like iron nuclei can be then efficiently accelerated in this scenario up to high energies ($\sim 10^{20}$ eV) by Fermi mechanism as we shall discuss in the next section.

III. CR ACCELERATION AT NGC 253

We suggest that the iron nuclei acceleration in NGC 253 occurs through a two-step process. In a first stage, ions are diffusively accelerated at single supernova shock waves within the nuclear region of the galaxy. Energies up to $\sim 10^{15}$ eV can be achieved in this step [23]. Fe-nuclei are not photodissociated in the process despite the starburst's central photon density is much larger than that of the Milky Way. The continuum spectrum of NGC peaks in the far infrared at $\sim 100\mu$m, with a luminosity of $\sim 3 \times 10^{10}$ $L_\odot$ [24] and a photon energy density of $U_{\rm ph} \sim 200$ eV cm^{-3} [13]. For such values is straightforward that the nuclei interactions with the blue–shifted ambient photons are quite below the photodisintegration threshold. Energy losses through pair production are also negligible.

By other hand, interactions with the rich interstellar gas in the center of the galaxy might disintegrate the ions if the escape from the starburst region is dominated by diffussion. Multiwavelength molecular line observations of the central region show that the average density of the molecular clouds is $\sim 10^5$ cm^{-3}, with a filling factor $< 10^{-3}$ [25]. This means an average gas density (mainly H_2) in the active region within the range 30 - 300 cm^{-3} [17]. The cross section for Fe$-$H$_2$ interactions ($\sigma \approx 1463.7$ mb) can be estimated from Rudstam's parametrization of proton induced spallation, $\sigma = 50 A_t^{2/3}$ mb (where A_t is the mass number of the target nucleus) [26]. The mean free path for the iron nuclei is then $\lambda \sim (n\sigma_{\rm Fe-H_2})^{-1}$ which results in the range 738 - 7380 pc when the upper and lower limits of particle density are considered. The central starburst region can be modeled as a disk of 70 pc thick with a radius $R \approx 300$ pc [17]. Since the gyroradius of the Fe nuclei is $\sim 10^{-3}$ pc, they certainly cannot be driven out from the starburst by diffusion.

However, due to the nature of the central region in NGC 253, the escape of the iron nuclei is expected to be dominated by convection. In fact, the presence of several tens of young SNRs with very high expansion velocities (~ 12000 km s^{-1} [27]) and thousands of massive O stars (with stellar winds of terminal velocities up to 3000 km s^{-1}) must generate collective plasma motions of several thousands of km per second. Then, due to the coupling of the magnetic field to the hot plasma, the magnetic field is also lifted outwards and forces the cosmic ray gas to stream along from the starburst region.

The relative importance of convection and diffusion in the escape of the cosmic rays from a region of disk scale height h is given by the dimensionless parameter,

$$q = \frac{V_0\, h}{\kappa_0}, \tag{1}$$

where V_0 is the convection velocity and κ_0 is the cosmic ray diffusion coefficient inside the starburst [28]. When $q \lesssim 1$, the cosmic ray outflow is "difussion dominated", whereas when $q \gtrsim 1$ it is "convection dominated". Assuming for the central region of NGC 253 a convection velocity of the order of the expanding SNR shells (i.e. $\sim$ 10000 km s^{-1}, a scale height $h \sim 35$ pc, and a reasonable value for the diffusion coefficient $\kappa_0 \sim 5 \times 10^{26}$ cm^2 s^{-1} [29], we get $q \sim 216$ and convection dominates the escape of the particles. The residence time of the iron nuclei in the starburst results $t_{\rm RES} \sim h/V_0 \approx 1 \times 10^{11}$ s. Most of the nuclei then escape through the disk in opposite directions along the symmetry axis of the system, being the total path traveled substantially shorter than the mean free path.

Once the nuclei escape from the central region of the galaxy with energies of $\sim 10^{15}$ eV, they are injected into the galactic-scale wind and experience further acceleration at its terminal shock.* The scale length of this shock is of the order of several tens of kpc (see Ref. [13]), so it can be considered as locally plane for calculations. The shock velocity $v_{\rm sh}$ can be estimated from the empirically determined superwind kinetic energy flux $\dot{E}_{\rm sw}$ and the mass flux $\dot{M}$ generated by the starburst through:

$$\dot{E}_{\rm sw} = \frac{1}{2}\dot{M}v_{\rm sh}^2. \tag{2}$$

The shock radius can be approximated by $r \approx v_{\rm sh}\tau$, where τ is the starburst age. Since the age is about a few tens of million years, the maximum energy attainable in this configuration is constrained by the limited acceleration time arisen from the finite shock's lifetime. For this second step in the acceleration process, the photon field energy density drops to values of the order of the cosmic background radiation (we are now far from the starburst region), and consequently, as we shall discuss in the next section, iron nuclei are again safe from photodissociation while energy increases from $\sim 10^{15}$ to 10^{20} eV.

In order to estimate the maximum energy that can be reached by the nuclei in the second stage of the acceleration process let us consider the superwind terminal shock propagating in a homogeneous medium with an average magnetic field B. If we work in the frame where the shock is at rest, the upstream flow velocity will be $\mathbf{v_1}$ ($|\mathbf{v_1}| = v_{\rm sh}$) and the downstream velocity, $\mathbf{v_2}$. The magnetic field turbulence is assumed to lead to isotropization and consequent diffusion of energetic particles which then propagate according to the standard transport theory [31]. The acceleration time scale is then [32]:

$$t_{\rm acc} = \frac{4\kappa}{v_1^2} \tag{3}$$

where κ is the upstream diffusion coefficient which can be written in terms of perpendicular and parallel components to the magnetic field, and the angle θ between the (upstream) magnetic field and the direction of the shock propagation:

$$\kappa = \kappa_\parallel \cos^2\theta + \kappa_\perp \sin^2\theta \tag{4}$$

Since strong turbulence is expected from the shock we can take the Bohm limit for the upstream diffusion coefficient parallel to the field, i.e.

$$\kappa_\parallel = \frac{1}{3}\frac{E}{ZeB_1} \tag{5}$$

where B_1 is the strength of the pre-shock magnetic field and E is the energy of the Z-ion. For the $\kappa_\perp$ component we shall assume, following Biermann [33], that the free mean path perpendicular to the magnetic field is independent of the energy and has the scale of the thickness of the shocked layer ($r/3$). Then,

$$\kappa_\perp = \frac{1}{3}r(v_1 - v_2) \tag{6}$$

or, in the strong shock limit,

$$\kappa_\perp = \frac{rv_1^2}{12}. \tag{7}$$

Since the upstream time scale is $t_{\rm acc} \sim r/(3v_1)$, we rewrite Eq. (3) as:

$$\frac{r}{3v_1} = \frac{4}{v_1^2}\left(\frac{E}{3ZeB_1}\cos^2\theta + \frac{rv_1^2}{12}\sin^2\theta\right), \tag{8}$$

and then, using $r = v_1\tau$ and transforming to the observer's frame, we get:

$$E_{\rm max} \approx \frac{1}{4}ZeBv_{\rm sh}^2\tau, \tag{9}$$

or

$$E_{\rm max} \approx \frac{1}{2}ZeB\frac{\dot{E}_{\rm sw}}{\dot{M}}\tau, \tag{10}$$

in terms of parameters that can be determined from observations.

The predicted kinetic energy and mass fluxes of the starburst of NGC 253 derived from the measured IR luminosity are 2×10^{42} erg s^{-1} and 1.2 M$_\odot$ yr^{-1}, respectively [13]. The starburst age is estimated from numerical models that use theoretical evolutionary tracks for individual stars and make sums over the entire stellar population at each time in order to produce the galaxy luminosity as a function of time [34]. Fitting the observational data these models provide a range of suitable ages for the starburst phase that, in the case of NGC 253, goes from 5×10^7 to 1.6×10^8 yr (also valid for M82) [34]. These models must assume a given initial mass function (IMF), which usually is taken to be a power-law with a variety of slopes. Recent studies has shown that the same IMF can account for the properties of both NGC 253 and M82 [21]. Here we shall assume a conservative age $\tau = 50$ Myr. Finally, the radio and γ-ray emission from NGC 253 are well matched by models with $B \sim 50\mu$ G [17]. With these figures we obtain a maximum energy for iron nuclei of:

$$E_{\rm max}^{\rm Fe} \sim 3.4\times10^{20} \quad \text{eV}, \tag{11}$$

a value quite lower (more than two orders of magnitude) than the limit imposed by the synchrotron losses [35].

*CR acceleration at superwind shocks was firstly proposed by Jokipii and Morfill during the 80s [30] in the context of our own Galaxy.

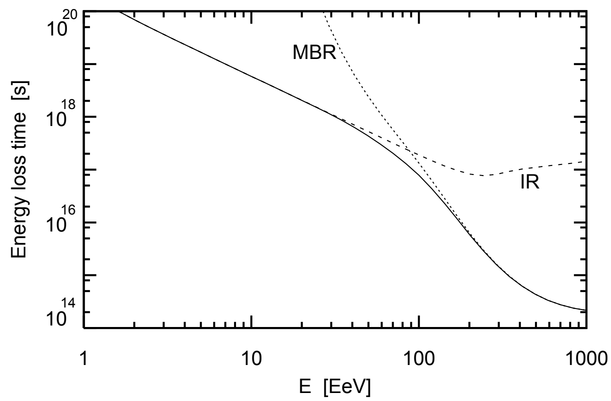

FIG. 1. Effective energy loss time for iron nuclei photodisintegration through the 2.73K and IR photons. Adapted from Epele and Roulet Ref. [7].

Regarding the spectral slope of the particles, it is expected that they emerge from the central region of NGC 253, where there are strong stellar winds from the massive star population, with an energy index $\gamma = 2.4$ [33], which is in fine accordance with the observed radio spectrum $\propto \nu^{-0.7}$ observed in the individual SNRs within the starburst [19]. After the particles left the nuclear region, diffusive losses tend to steepen the spectrum at the terminal shock of the galactic superwind, where reacceleration takes place. The final index will depend on some not well known parameters as the Mach number of the terminal shock. In what follows, we shall study the propagation of the nuclei from NGC 253 to the Earth in models with $\gamma = 2.4$, 2.5, and 2.6.

IV. PROPAGATION EFFECTS

The basic interactions between the universal background radiation and nuclei of extremely high energy have been first discussed in detail by Puget *et al* [36]. Heavy nuclei with energies above a few EeV get attenuated mainly by photodisintegration off the microwave background radiation (MBR) and the intergalactic infrared (IR) background photons.†

The photodisintegration process is dominated by a broad maximum designated as the giant dipole resonance which peaks in the γ-ray energy range of 10 to 30 MeV (nucleus rest frame). In the initial absorption process, the photon energy may be given to a single nucleon (although almost always is shared with others), so the decay process might involve the emission of one (or more) nucleons. The disintegration rate of a nucleus of mass A with the subsequent production of i nucleons reads [37],

$$R_{Ai} = \frac{1}{2\Gamma^2}\int_0^\infty dw\,\frac{n(w)}{w^2}\int_0^{2\Gamma w} dw_r\, w_r \sigma_{Ai}(w_r) \qquad (12)$$

where $n(w)$ is the density of photons with energy w in the system of reference in which the MBR is at 2.73 K. In the formula, w_r is the energy of the photons in the rest frame of the nucleus, Γ is the Lorentz factor, and σ_{Ai} is the cross section for the interaction. The total cross section of photonuclear interaction is well known [38]. As usual, the MBR is described by Plankian spectrum of 2.73 K. For the diffuse IR radio background (photon energies between 2×10^{-3} - 8×10^{-1} eV) we shall use the estimate given in Ref. [39],

$$\frac{dn}{dw} = 1.1\times10^{-4}\left(\frac{w}{\rm eV}\right)^{-2.5}\ {\rm cm^{-3}eV^{-1}}. \qquad (13)$$

With these figures at hand it is straightforward to compute the energy loss time of iron nuclei (displayed in Fig. 1). It is clear from these calculations that the universal background radiation would not affect iron acceleration in NGC 253 up to a few thousands of EeV. Moreover, since the nearness of the starburst galaxy, just the interaction with the MBR becomes relevant in their travel to Earth. The fractional energy loss as a function of the MBR energy (Lorentz factor) have been already parametrized [8],

$$R(\Gamma) = 3.25\times10^{-6}\,\Gamma^{-0.643}\exp(-2.15\times10^{10}/\Gamma)\ {\rm s}^{-1} \qquad (14a)$$

if $\Gamma \in [1.\times10^9, 3.68\times10^{10}]$, and

$$R(\Gamma) = 1.59\times10^{-12}\,\Gamma^{-0.0698}\ {\rm s}^{-1} \qquad (14b)$$

if $\Gamma \in [3.68\times10^{10}, 1.\times10^{11}]$. It is noteworthy that the possible disintegration histories computed using Eqs. (14a) and (14b) are in very good agreement with Monte Carlo simulations when the sources are located near the Earth (distances $\lesssim 10$ Mpc) [40].

Using the formalism sketched in [8], it is easily obtained the evolution of the differential spectrum $Q(E_g,t) = KE_g^{-\gamma}\delta(t-t_0)$ of the iron nuclei injected by NGC 253 at t_0. The number of surviving fragments with energy E at time t is given by,

$$N(E,t)dE = \frac{KE_g^{-\gamma+1}}{E}dE, \qquad (15)$$

where E_g denotes the energy at which the nuclei were emitted from the source, related to the energy detected on Earth by $E = E_g\ e^{-R(\Gamma)\,t/56}$ (recall that the Lorentz

†Actually, the pair creation process due to interactions with the MBR as well as disintegration with the optical background photons also attenuate the nuclei energy. However, these processes are not essential for the discussion presented in this paper. For details of these interactions the reader is referred to Ref. [5].

factor, $\Gamma = E_g/56$, does not result modified during the propagation). [‡]

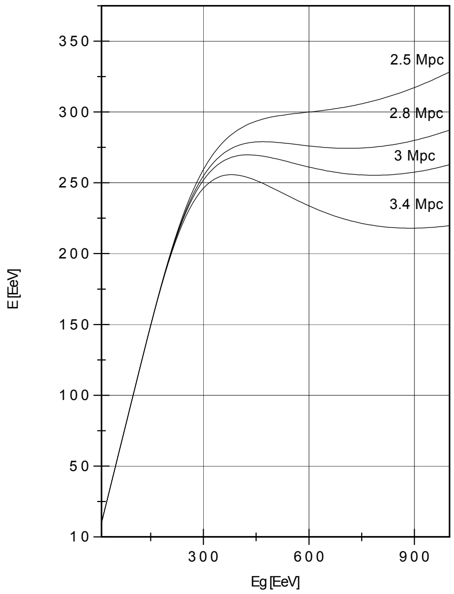

FIG. 2. Relation between the injection and arrival energies for possible propagation distances from NGC 253 according to current data.

In Fig. 2 it is shown the energy degradation for different flying timescales arising from different propagation regimes. The energy spectrum of the surviving fragments is degenerate. For instance, for a propagation distance of 3 Mpc, the composition of the arrival nuclei changes from $A = 55$ (for $\Gamma \approx 3 \times 10^9$) to $A = 44$ (for $\Gamma \approx 6 \times 10^9$). If a maximum energy of 340 EeV is attainable in the source, the relation between the injection and arrival energies is a monotonously decreasing function. At this stage, we conveniently introduce the modification factor η, defined as the ratio between the modified spectrum and the unmodified one. Notice that once the nuclei have energy enough to undergo photodisintegration through the giant dipole resonance, the value of the modification factor is always less than one.

Contrariwise, if the injection energy has an upper cutoff at $E \gtrsim 560$ EeV (this can be achieved, for instance, with a higher value of τ), the function which relates the injection and arrival energies becomes multivalued, yielding a bump-like feature in the modification factor. To understand this behavior, recall that nuclei suffer a violent disintegration at these energies (see Fig. 1), in such a way that, for a Lorentz factor $\Gamma \approx 8 \times 10^9$, the composition of the arrival nuclei drops to $A \approx 33$. So, particles injected with different energies might arrive with the same energy, piling up around 250 EeV, just before the expected cutoff. The whole effect can be clearly appreciated in Fig. 3, where we have plotted the modification factor for two different cutoffs in the injection spectrum.

For longer propagation distances, the pile–up shifts to lower energies (e.g. to ≈ 200 EeV for 3.4 Mpc), broadening its profile. On the other hand, for shorter distances to the nuclei–emitting source there is no spectral bump, i.e., the spectrum is monotonic. Thus, the probability of detecting events diminishes with increasing Lorentz factor.

We remark that the change of the spectral index does not significantly affect the shape of the modification factor. The height of the pile–up diminishes with the steepening of the injected spectrum, but this effect is accompanied by a simultaneous drop in the CR flux of the preceding bin of energy, resulting in the same overall shape for η, although downshifted (see Figs. 4 and 5 of ref [8]).

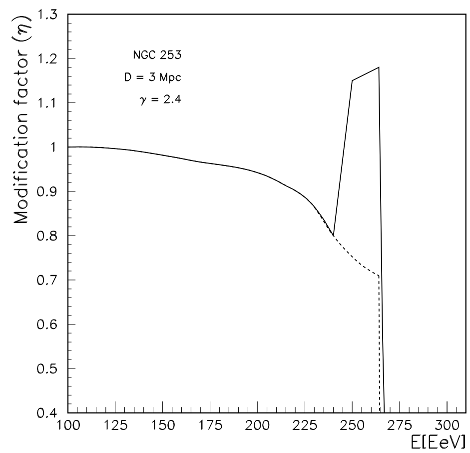

FIG. 3. Modification factor for initial iron nuclei from NGC 253, assuming a differential power law injection spectrum with spectral index $\gamma = 2.4$. Solid (dotted) line stands for the case with an upper cutoff at the injection energy of 560 EeV (340 EeV).

Events from the pile–up are about 50% more probable to be detected than those at energies immediately lower. The reason is that nuclei with energies $\gtrsim 560$ EeV would be almost completely photodisintegrated during their journey to Earth, in such a way that the surviving fragments end at energies of the pile–up, changing the relative detection probabilities. Thus, we have the interesting result that the existence of a sufficient high cutoff (at Lorentz factors, say, $\Gamma \gtrsim 8 \times 10^9$) in the source acceleration mechanism favors the detection of events around 250 EeV from the starburst galaxy NGC 253.

V. CONCLUDING REMARK

Might heavy nuclei be primaries at the end of the cosmic ray spectrum? If the spectrum extends over the up-

[‡]Notice that the (Z, A)-dependence of R is roughly cancelled by dividing by A in the exponent. This implies that R can in fact be integrated down to lower, spallated A values and still be reasonably accurate [36].

to-now observed energies, the answer will be certainly "no". Heavy nuclei with energies above 200 EeV could not propagate for more than 10 Mpc. Besides, Lorentz factors above 2×10^{10} require large astrophysical regions for acceleration where the relic photon density would be sufficient to provide a significant loss mechanism for the nuclei (see Fig. 1).

However, if the highest cosmic ray energy is of the order of 300 EeV, heavy nuclei accelerated in the nearby starburst galaxies (NGC 253 and M82) can be present at the end of the spectrum. We have shown that these nuclei, originated in the central regions of the galaxies, can be accelerated without suffering catastrophic interactions in a two-step process that involves supernova remnant shock waves, and the large scale terminal shock produced by the superwind that flows from the starbursts.

Anchordoqui *et al.* [8] have suggested that the lack of data at energies immediately lower than the two "super–GZK events" recorded to date, could be the result of a different primary composition of disintegrated iron nuclei at the end of the spectrum. This speculation might find some support from our calculations since they suggest that the production of the energetic nuclei in next-door galaxies (NGC 253 and M82) is at least feasible. In any case, we have made some concrete predictions that could be tested with the forthcoming facilities of the Southern Auger Observatory [41], and the satellite experiment called OWL [42].

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