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Cosmic ray photodisintegration and the knee of the spectrum

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Abstract

We explore in some detail the scenario proposed to explain the observed knee of the cosmic ray (CR) spectrum as due to the effects of photodisintegration of the CR nuclei by interactions with optical and soft UV photons in the source region. We show that the photon column densities needed to explain the experimental data are significantly lower than those obtained in previous estimations which neglected multinucleon emission in the photodisintegration process. We also treat more accurately the photodisintegration thresholds, we discuss the effects of photopion production processes and the neutron escape mechanism, identifying the physical processes responsible for the qualitative features of the results. This scenario would require the CR nuclei to traverse column densities of $\sim 5 \times 10^{27} - 2 \times 10^{28}$ eV/cm² after being accelerated in order to reproduce the observed knee, and predicts that the CR composition should become lighter above $\sim 10^{16}$ eV. © 2002 Elsevier Science B.V. All rights reserved.

1. Introduction

It is widely accepted nowadays that cosmic rays (CRs) with energies per particle up to about 10^{18} eV are protons and nuclei of galactic origin. Furthermore, it is well established that the full CR energy spectrum has a power-law behavior with a steepening taking place at the so-called knee, corresponding to an energy $E_{\text{knee}} = 3 \times 10^{15}$ eV. Although it is well known that the CR composition below the knee has a significant heavy component, its behavior beyond the knee remains somewhat controversial. Indeed, above 10^{14} eV/nucleus the information about the CR mass composition has to be drawn from extensive air shower observa-

tions at ground level with the results from different experiments not always being compatible. Moreover, the theoretical predictions depend on the hadronic model adopted and this can introduce further uncertainties on the inferred composition. The model dependence in certain observables can however be useful to discriminate among the hadronic models using the experimental data, and in this comparison QGSJET and VENUS turn indeed to be favored against other models such as SIBYLL [1]. Thus, the CR mass composition beyond the knee is not definitively established yet, with some observations [2] suggesting that it turns lighter and others instead suggesting that the heavier components become dominant [3–7].

The attempts made so far to explain the physical origin of the knee of the cosmic ray spectrum can be roughly classified into three kind of models. One of them exploits the possibility that the

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acceleration mechanisms could be less effective above the knee [8–10], while a second one assumes that leakage from the Galaxy plays the dominant role in suppressing CRs above the knee [11–13]. The third scenario, originally proposed by Hillas [14], considers nuclear photodisintegration processes (and proton energy losses by photomeson production), in the presence of a background of optical and soft UV photons in the source region, as the main responsible for the change in the CR spectrum. The first two scenarios are based on a rigidity dependent effect which consequently produces a change in the spectral slope of each nuclear component with charge Z at an energy $\sim E_{\text{knee}}Z$. Hence, it removes first the protons and the lighter nuclei and only at larger energies affects the heavier nuclei, predicting that the CR mass composition should become heavier beyond the knee. On the contrary, the third scenario predicts a lighter CR composition above the knee due to the disintegration of the propagating nuclei. The predictions of this last scenario have been worked out by Karakula and Tkaczyk [15] some years ago. The aim of the present paper is to re-evaluate in detail the propagation of nuclei and protons in the source region following their approach, but taking into account some additional features that turn out to be important for the determination of the source parameters required for the scenario to work. In particular, we re-evaluate the nuclear photodisintegration rates taking into account multinucleon emission processes, we introduce more accurate threshold energies in the giant dipole resonance for every stable nucleus (according to recent remarks by Stecker and Salamon [16]), and discuss the impact of the neutron mechanism [17,18] which would allow the escape of neutrons from the source without further energy losses. Moreover, we show that by introducing a non-negligible lower cutoff to the power-law photon distribution, the abrupt suppression of the all-particle CR spectrum above 10^{17} eV previously obtained [15] is prevented, and hence the underlying reason for that effect is identified. Let us also mention that although we focus our study here on the scenario proposed to explain the observed steepening at the knee, we expect that our general considerations about the treatment of photodis-

integration processes should also be relevant in other contexts.

2. The propagation of cosmic rays

2.1. The source and the photon background spectra

As already mentioned, we deal with a model that considers that the CR nuclei are accelerated inside discrete sources which are surrounded by a very strong background of optical and soft UV photons. If the typical energy of the photons is in the optical range (1–10 eV), the photodisintegration of CR nuclei will start to be efficient at CR energies $E \geq A \times 10^{15}$ eV. This is so because in the CR rest frame the optical photon (which is boosted by a relativistic factor $\gamma \geq 10^6$) appears as an energetic gamma ray with $E \geq 1$ –10 MeV, i.e., capable of photodisintegrating the nucleus. Hence, the basic assumption is that the CR spectrum emitted by the source is a featureless extrapolation of the spectra measured below the knee and that the observed change in slope is actually due to the CR interactions (i.e., the process of nuclear photodisintegration and the photopion production by protons) with the surrounding photon background.

We then assume that the source emits nuclei with mass number A (with A ranging from 1 to 56). Since there is only one stable isotope for a nucleus of a given mass A along most of the decay chain from ^{56}Fe to ^1H , we take for definiteness a unique charge Z associated to any given value of A . The differential fluxes emitted by the source are assumed to be given by power-law distributions

$$\phi_i^0(E) = \Phi_i^0 E^{-\gamma_i} \quad (1)$$

(where $i = 56 - A$ throughout). The intensities Φ_i^0 and spectral indices γ_i were taken from the detailed knowledge about CR mass composition below the knee (for energies per particle above $\sim \text{few } Z \times 10^{10}$ eV) [19].

For the photon background that surrounds the source we consider that its spectrum follows either a (thermal) Planckian-type distribution given by

$$n(\epsilon) = \frac{\rho}{2\zeta(3)(k_B T)^3} \frac{\epsilon^2}{\exp(\epsilon/k_B T) - 1}, \quad (2)$$

where k_B is the Boltzmann constant, T the absolute temperature and ζ the Riemann zeta function (such that $\zeta(3) \approx 1.202$), or a power-law distribution given by

$$n(\epsilon) = \rho(\alpha - 1) \left(\frac{1}{\epsilon_m^{\alpha-1}} - \frac{1}{\epsilon_M^{\alpha-1}} \right)^{-1} \epsilon^{-\alpha}, \quad (3)$$

where ϵ_m, ϵ_M are the lower and upper energy cut-offs and α the assumed spectral index. For both distributions, ρ represents the total number density and $n(\epsilon)$ the differential number density corresponding to a photon energy ϵ . Defining ρ_E as the total energy density, one has for the Planckian spectrum

$$\rho_E \approx 2.631 k_B T \rho, \quad (4)$$

while the analogous expression for the power-law spectrum is

$$\rho_E = \frac{\alpha - 1}{2 - \alpha} (\epsilon_M^{2-\alpha} - \epsilon_m^{2-\alpha}) \left(\frac{1}{\epsilon_m^{\alpha-1}} - \frac{1}{\epsilon_M^{\alpha-1}} \right)^{-1} \rho. \quad (5)$$

We adopt $\alpha = 1.3$ as in Ref. [15], but the introduction of a non-negligible lower cutoff ϵ_m in the power-law spectrum excludes the presence of an abundant number of low energy (infrared) photons that would otherwise play the dominant role in the photodisintegration above 10^{17} eV, as was the case in Ref. [15]. Regarding the spatial dependence of the radiation background we do not make any particular assumption, but the results will depend only on the integrated photon column density along the trajectory of the CR since its production at the source until it leaves the region filled with radiation.

2.2. Propagation of nuclei: photodisintegration

The main mechanism of energy loss for nuclei with energy per particle below 10^{18} eV propagating through these photon backgrounds is the process of photodisintegration. In fact, a nucleus of mass $A = 56 - i$ and Lorentz factor $\gamma = E/Am_p c^2$ that propagates through a photon distribution $n(\epsilon)$ (as,

for example, those of Eqs. (2) and (3)) has a probability of fission with emission of j nucleons given by

$$R_{ij}(E) = \frac{1}{2\gamma^2} \int_{\epsilon'_{\text{thr},ij}/2\gamma}^{\infty} d\epsilon \frac{n(\epsilon)}{\epsilon^2} \int_{\epsilon'_{\text{thr},ij}}^{2\gamma\epsilon} d\epsilon' \epsilon' \sigma_{ij}(\epsilon'), \quad (6)$$

where σ_{ij} is the corresponding photodisintegration cross-section, ϵ the photon energy in the observer's system and ϵ' its energy in the rest frame of the nucleus. In order to calculate R_{ij} we fitted σ_{ij} with the parameters given in Tables 1 and 2 of Ref. [20], while the reaction thresholds $\epsilon'_{\text{thr},ij}$ were taken from Table 1 of Ref. [16]. A useful quantity to take into account all reaction channels is the effective emission rate given by

$$R_{i,\text{eff}} = \sum_{j \geq 1} j R_{ij}. \quad (7)$$

At low energies ($\epsilon'_{\text{thr},ij} \leq \epsilon' \leq 30$ MeV), the cross-section $\sigma_{ij}(\epsilon')$ is dominated by the giant dipole resonance and the photodisintegration proceeds chiefly by the emission of one or two nucleons. Note that, according to recent remarks pointed out in Ref. [16], we employ threshold energies that depend both on the nucleus and on the number of emitted nucleons. The improvement consists in shifting the threshold energy, that was previously assumed to be $\epsilon'_{\text{thr}} = 2$ MeV for all reaction channels and all nuclei [15,20], to higher energies, such that single-nucleon emission has now a typical threshold of ~ 10 MeV, while the double-nucleon emission energy threshold becomes typically ~ 20 MeV. Although the appearance of the knee in this scenario is ultimately a threshold effect, the improved calculation including the change in the specific threshold energies has however a minor impact on the resulting emission rates when compared with those obtained with a fixed threshold energy $\epsilon'_{\text{thr}} = 2$ MeV. This can be appreciated in Fig. 1, where plots of $R_{i,\text{eff}}/\rho_E$ vs energy per nucleon E^* are shown for a ^{56}Fe nucleus ($i = 0$) propagating through a thermal photon spectrum (with $k_B T = 10$ eV). The emission rate $R_{0,\text{eff}}$ (calculated with the specific threshold energies) shows a small shift towards higher energies with respect to that obtained by means of the fixed threshold energy (labeled $R_{0,\text{eff}}^{(2)}$). The reason for this reduced

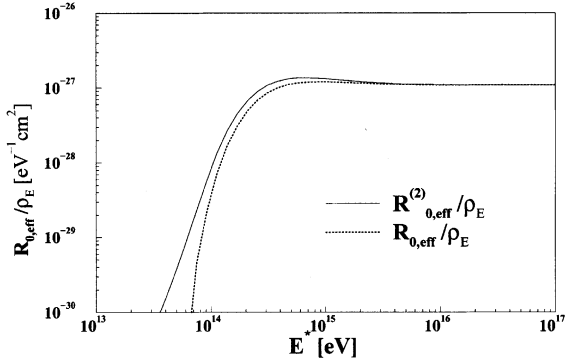


Fig. 1. Plots of $R_{0,\text{eff}}/\rho_E$ versus energy per nucleon E^* for a thermal photon spectrum (with $k_B T = 10$ eV). $R_{0,\text{eff}}^{(2)}$ was evaluated using the fixed threshold energy $\epsilon_m = 2$ MeV, while $R_{0,\text{eff}}$ corresponds to the improved threshold energies according to Ref. [16].

effect of the threshold shifts is that the dominant photodisintegration effects are due anyhow to the energies around the giant resonance peak, which for single-nucleon emission are typically peaked around ~ 20 MeV and for double-nucleon emission around ~ 26 MeV, and hence clearly well above the threshold energies.

At higher energies, $\sigma_{ij}(\epsilon')$ is approximately flat and the multinucleon emission acquires a higher probability, becoming actually dominant for heavy nuclei, for which the probability for single nucleon emission is of only 10%. To estimate the relevance of multinucleon emission processes, we compare the effective rate from Eqs. (6) and (7) with the emission rate that neglects multinucleon processes (as calculated in Ref. [15]). Fig. 2(a) shows plots of $R_{i,\text{eff}}/\rho_E$ versus energy per nucleon E^* for a ^{56}Fe nucleus ($i = 0$) propagating through a thermal photon spectrum (with $k_B T = 1.8$ eV). From the figure it becomes evident that multinucleon processes play an important role, increasing the emission rate at high energies by a factor of ~ 4 with respect to the single nucleon emission results. Analogously, Fig. 2(b) shows the corresponding plots for a power-law photon spectrum with upper energy cutoff $\epsilon_M = 20$ eV and spectral index $\alpha = 1.3$. From this figure we can also see the effect of introducing a non-negligible lower cutoff in the spectrum. In fact, we observe that with a negligible cutoff (e.g., $\epsilon_m = 10^{-6}$ eV) the emission rate grows

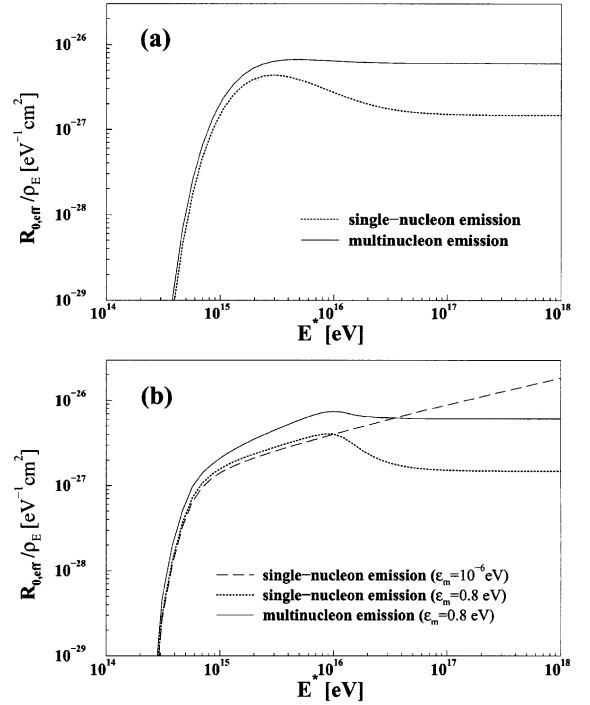


Fig. 2. Plots of $R_{0,\text{eff}}/\rho_E$ versus energy per nucleon E^* . (a) Results for a thermal photon spectrum with $k_B T = 1.8$ eV. (b) Results for a power-law photon spectrum with upper energy cutoff $\epsilon_M = 20$ eV and spectral index $\alpha = 1.3$. The multinucleon emission rate corresponds to a lower energy cutoff $\epsilon_m = 0.8$ eV, while the single-nucleon emission rates correspond to $\epsilon_m = 10^{-6}$, 0.8 eV, as indicated.

steadily, because increasingly abundant low energy (infrared) photons give in this case the dominant contribution to the rates. They also keep the influence of the giant resonance dominant even at high particle energies, and hence in this case the multinucleon emission has little impact. However, if we take a non-negligible cutoff into account (for instance, $\epsilon_m = 0.8$ eV), we find that the flat high energy regime of σ_i now dominates beyond a given value of E^* and hence the emission rate saturates at high energies. Comparing the two emission rate curves with non-negligible cutoff, we observe again that multinucleon emission processes lead to a significant increase (by the same factor of ~ 4 we already encountered) in the emission rate. Thus, introducing a non-negligible lower cutoff energy in the power-law photon spectrum the results should

not differ much from those with the thermal photon spectrum and, being the rates with multinucleon emission larger, they should also require smaller photon densities than those previously estimated in Ref. [15] to produce similar overall effects.

It should be noticed that in Figs. 1 and 2 we have chosen to plot the emission rates as functions of the energy per nucleon E^* , instead of energy per nucleus E , because it is E^* that remains constant during photodisintegration and also because for different nuclei the shapes of the cross-section are similar with the maxima occurring approximately at the same values of E^* .

The equations that describe the propagation of nuclei are:

$$\frac{\partial \phi_i(E^*, x)}{\partial x} = -\phi_i(E^*, x) \sum_{j \geq 1} R_{ij}(E^*) + (1 - \delta_{0i}) \times \sum_{j=1}^i R_{(i-j)j}(E^*) \phi_{i-j}(E^*, x), \quad (8)$$

where $\phi_i(E^*, x)$ is the differential flux of a nucleus of mass $A = 56 - i$ (for $0 \leq i \leq 54$) with an energy per nucleon E^* at propagation distance x from the source, $R_{ij}(E^*)$ is the above defined emission rate and δ_{ij} is the Kronecker delta. The exact solution of these equations (i.e., the differential fluxes at propagation distance L) is given by

$$\phi_i(E^*, L) = \sum_{j=0}^i b_{ij}(E^*) \exp\left(-\sum_{k \geq 1} R_{jk}(E^*)L\right), \quad (9)$$

where

$$b_{ii} = \phi_i^0(E^*) - (1 - \delta_{0i}) \sum_{j=0}^{i-1} b_{ij}(E^*) \quad (10)$$

and

$$b_{ij} = \frac{\sum_{k=1}^{i-j} b_{(i-k)j}(E^*) R_{(i-k)k}(E^*)}{\sum_{k \geq 1} R_{ik}(E^*) - \sum_{k \geq 1} R_{jk}(E^*)} \quad \text{for } i > j. \quad (11)$$

Recalling that the emission rates R_{ij} are linear in the photon energy density ρ_E , it should be noted that, as expected, the fluxes only depend on the column density $\rho_E L$. It also should be pointed out that, since the equations describing the CR prop-

agation only depend on the energies per nucleon E^* , the final CR spectra require the re-calculation of the solution (Eqs. (9)–(11)) in terms of energy per nucleus.

2.3. Propagation of protons

The main energy loss mechanism that must be taken into account for the propagation of protons is the photomeson production process. We have to distinguish between the primary protons produced at the source with an initial flux $\phi_{i=55}^0$ (as given by Eq. (1)) and the nucleons released as by-products of the nuclear photodisintegration. The fluxes corresponding to these (secondary) protons and neutrons (ϕ^{sp} and ϕ^{n} , respectively) can be easily determined from the solution for the nuclear fluxes, as we will see in the following. Let $\phi_i^{\text{N}}(E, x)$ be the differential flux of nucleons of energy E produced within a distance x away from the source by the photodisintegration of nuclei of mass number $A = 56 - i$ and energy per nucleon $E^* = E/A$. Then, the equation governing the evolution of the fluxes is

$$\frac{\partial \phi_i^{\text{N}}(E, x)}{\partial x} = \phi_i(E^*, x) \sum_{j \geq 1} j R_{ij}(E^*) \quad \text{for } 0 \leq i \leq 54 \quad (12)$$

(with the initial condition $\phi_i^{\text{N}}(E, x=0) = 0$) and its solution reads

$$\phi_i^{\text{N}}(E, L) = \sum_{j=0}^i \left(\frac{\sum_{k \geq 1} k R_{ik}(E^*)}{\sum_{k \geq 1} R_{jk}(E^*)} \right) b_{ij}(E^*) \times \left[1 - \exp\left(-\sum_{k \geq 1} R_{jk}(E^*)L\right) \right]. \quad (13)$$

Thus, the total fluxes of secondary protons and neutrons are found by performing the weighted sums over index i :

$$\phi^{\text{sp}}(E, L) = \sum_{i=0}^{54} \frac{Z_i}{56-i} \phi_i^{\text{N}}(E, L), \quad (14)$$

$$\phi^{\text{n}}(E, L) = \sum_{i=0}^{54} \left(1 - \frac{Z_i}{56-i} \right) \phi_i^{\text{N}}(E, L), \quad (15)$$

where Z_i is the charge of the i -nucleus.

We will also discuss the possible impact of the so-called neutron mechanism [17,18], which could take place in the presence of a magnetic field in the source region. Indeed, a magnetic field would retain the charged particles within that region for some period of time, while non-charged particles would be able to escape rapidly. In that case, we will consider that the flux of neutrons ϕ^n adds to the final all-particle CR flux neglecting any energy losses, while the flux of secondary protons ϕ^{sp} , combined with the initial flux of primary protons $\phi_{i=55}^0$, will undergo energy losses due to the well studied process of photopion production. Indeed, since a background photon looks like a high-energy gamma ray in the proton rest frame, the photopion reactions



proceed with an interaction probability given by

$$g(E) = \frac{1}{2\gamma^2} \int_{\epsilon'_{thr}/2\gamma}^{\infty} d\epsilon \frac{n(\epsilon)}{\epsilon^2} \int_{\epsilon'_{thr}}^{2\gamma\epsilon} d\epsilon' \epsilon' \sigma_{\gamma p}(\epsilon') K(\epsilon'), \quad (18)$$

where, analogously to Eq. (6), $\gamma = E/m_p c^2$ is the Lorentz factor of the proton, $\sigma_{\gamma p}$ is the γp interaction cross-section (parametrized from experimental data compiled in Ref. [21]), K is the coefficient of inelasticity (defined as the average relative energy loss of the proton), ϵ the photon energy in the observer's system and ϵ' its energy in the rest frame of the proton. The energy threshold for photopion production is $\epsilon'_{thr} = 145$ MeV.

Then, if we call $\phi^p(E, x)$ and $\phi^{*n}(E, x)$ the differential fluxes of protons and neutrons of energy E (produced by the reactions given by Eqs. (16) and (17), respectively) after traversing a distance x from the source, the corresponding propagation equations are

$$\begin{aligned} \frac{\partial \phi^p(E, x)}{\partial x} &= -g(E)\phi^p(E, x) \\ &+ \frac{1}{2} \frac{1}{(1-K)} g\left(\frac{E}{1-K}\right) \phi^p\left(\frac{E}{1-K}, x\right) \end{aligned} \quad (19)$$

and (neglecting neutron energy losses)

$$\frac{\partial \phi^{*n}(E, x)}{\partial x} = \frac{1}{2} \frac{1}{(1-K)} g\left(\frac{E}{1-K}\right) \phi^p\left(\frac{E}{1-K}, x\right), \quad (20)$$

where the factors 1/2 arise from the relative probability of occurrence of reactions (16) and (17). The coefficient of inelasticity K was given a typical value $K = 0.3$, irrespective of proton energy [15]. The initial conditions to solve these differential equations are

$$\phi^p(E, x=0) = \phi_{i=55}^0(E) + \phi^{sp}(E, L) \quad (21)$$

and

$$\phi^{*n}(E, x=0) = 0, \quad (22)$$

respectively. In the first one we assumed for simplicity that the secondary protons produced by photodisintegration of nuclei add directly to the initial proton spectrum. This is justified since whenever the photopion processes are relevant, the photodisintegration is very efficient and hence nuclei are rapidly disintegrated. The solution to Eq. (19) turns out to be

$$\phi^p(E, L) = \sum_{l=0}^{\infty} \phi_l^p(E, L), \quad (23)$$

where

$$\phi_0^p(E, L) = \phi^p(E, 0) \exp(-g(E)L), \quad (24)$$

while for $l > 0$

$$\begin{aligned} \phi_l^p(E, L) &= \frac{\phi^p\left(\frac{E}{(1-K)^l}, 0\right)}{(2(1-K))^l} \prod_{j=1}^l g\left(\frac{E}{(1-K)^j}\right) \\ &\times \sum_{n=0}^l \frac{\exp\left(-g\left(\frac{E}{(1-K)^n}\right)L\right)}{\prod_{m=0, m \neq n}^l \left[g\left(\frac{E}{(1-K)^m}\right) - g\left(\frac{E}{(1-K)^n}\right)\right]}. \end{aligned} \quad (25)$$

For the neutron component, the solution to Eq. (20) reads

$$\phi^{*n}(E, x) = \sum_{l=0}^{\infty} \phi_l^{*n}(E, x), \quad (26)$$

where

$$\phi_0^{*n}(E, L) = \frac{\phi^p\left(\frac{E}{1-K}, 0\right)}{2(1-K)} \left(1 - \exp\left(-g\left(\frac{E}{1-K}\right)L\right)\right), \quad (27)$$

while for $l > 0$

$$\begin{aligned} \phi_l^{*n}(E, L) &= \frac{\phi^p\left(\frac{E}{(1-K)^{l+1}}, 0\right)}{(2(1-K))^{l+1}} \prod_{j=0}^l g\left(\frac{E}{(1-K)^{j+1}}\right) \\ &\times \sum_{n=0}^l \frac{1 - \exp\left(-g\left(\frac{E}{(1-K)^{n+1}}\right)L\right)}{g\left(\frac{E}{(1-K)^{n+1}}\right) \prod_{m=0}^l \left[g\left(\frac{E}{(1-K)^{m+1}}\right) - g\left(\frac{E}{(1-K)^{n+1}}\right)\right]}. \end{aligned} \quad (28)$$

The linearity of the photopion reaction rate $g(E)$ with respect to the photon energy density ρ_E (as implied by Eq. (18)) causes the proton and neutron fluxes to depend only on the column density $\rho_E L$, similarly as for the fluxes of nuclei.

3. Results and comparison with observations

So far, we wrote down the appropriate equations governing the propagation of CR particles, and we found their exact solutions: Eqs. (9)–(11) for nuclei, Eqs. (13), (15) and (26)–(28) for neutrons (that eventually decay into protons outside the source region), and Eqs. (23)–(25) for protons (that undergo energy losses due to photopion production), the sum of all these contributions being the total CR flux ϕ_{total} . We can now investigate how ϕ_{total} depends on the assumed photon distributions around the source, namely on the column density $\rho_E L$ and on the parameters involved in the photon spectral distributions (Eqs. (2) and (3)).

Fig. 3(a) shows the total CR differential flux as a function of the energy per particle for a Planckian spectrum with $k_B T = 1.8$ eV and several values for the column density $\rho_E L$. Analogously, Fig. 3(b) shows the results corresponding to a Planckian spectrum with $k_B T = 10$ eV. To compare our results with experimental data, the figures also show the observed spectra measured by differential experiments (CASABLANCA, DICE, Tibet, PROTON satellite and AKENO) [1,2,22].

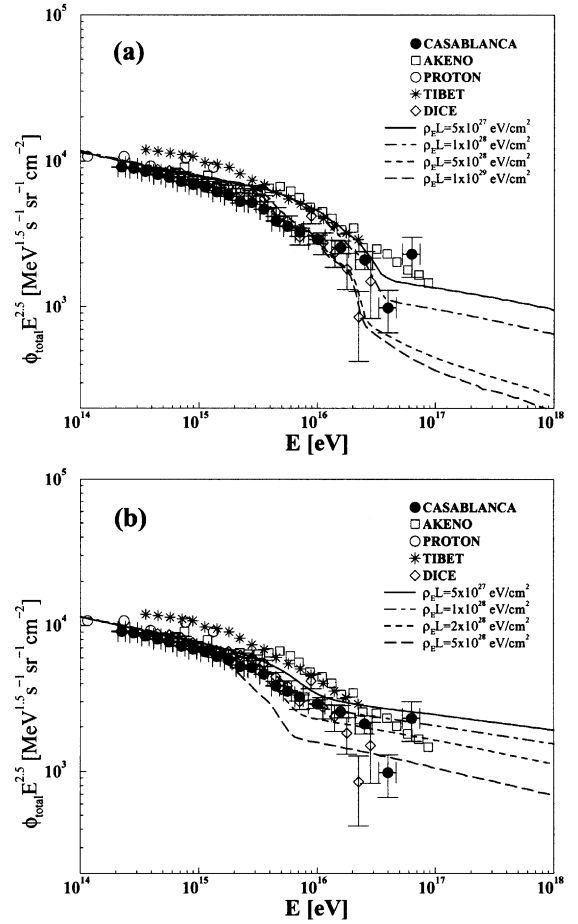


Fig. 3. Plots showing the total CR flux versus the energy per particle for thermal spectra with (a) $k_B T = 1.8$ eV and (b) $k_B T = 10$ eV, and several values for the column density $\rho_E L$, as indicated. For the sake of comparison, experimental data measured by different experiments are also shown.

The effect of increasing the temperature is clearly that of shifting the knee towards lower energies, while the effect of increasing $\rho_E L$ is that of intensifying the steepening at the knee. This behavior is easily explained by the fact that it is the nuclear photodisintegration the responsible for the occurrence of the knee (since photopion production processes show up at higher energies). Then, if the photon temperature (and thus, the mean photon energy) is larger, the particle energy required for the nuclei to disintegrate decreases. Furthermore, CR propagation cannot be affected

by $\rho_E L$ below the knee, where the photodisintegration has not set up yet, but above the knee the disintegration rates have a linear dependence on ρ_E . Hence, we should expect $\rho_E L$ to control the steepening in the knee region, in agreement with the results shown in the figures.

In Ref. [15] the best fit to observations for the Planckian spectrum was that corresponding to $k_B T = 1.8$ eV and $\rho_E L = 2.25 \times 10^{29}$ eV/cm². For the sake of comparison, Fig. 3(a) shows our results for the same temperature. It is clearly observed that the new results are compatible with the AKENO experimental data for low column density values $\rho_E L \simeq 5 \times 10^{27}$ eV/cm², i.e., ~ 50 times lower than the previous result. Fig. 3(b) shows that the corresponding results for $k_B T = 10$ eV are in somewhat better agreement with CASABLANCA observations for column densities in the range $\rho_E L = 5 \times 10^{27} - 2 \times 10^{28}$ eV/cm². One has to keep in mind however that the model discussed is oversimplified in many respects: we assume that the same photon column density is traversed by all nuclear species and for all CR energies, we neglect additional effects due to change in diffusion properties at these energies, which in alternative scenarios [13] are actually the responsible for the appearance of the knee, we assume that CR interactions in the interstellar medium are negligible, etc. Hence, an accurate agreement with observational data is not to be expected, although it is reassuring to see that the results have the right tendency to account for the main qualitative features of the data in the knee region.

Fig. 4(a) and (b) shows plots of mean mass composition $\langle \ln A \rangle$ versus E corresponding to the Planckian photon spectra with $k_B T = 1.8$ eV and $k_B T = 10$ eV, respectively. For comparison, we also show the CR mass composition measured by different experiments [1,2,5,6,23]. One can observe that solutions corresponding to the largest plotted values of column density $\rho_E L$ imply that CRs above the knee are chiefly constituted by protons, in disagreement with experimental data; the non-negligible contribution of heavier CR components beyond the knee imposes an upper limit on the column density around $(\rho_E L)_{\max} \approx 2 \times 10^{28}$ eV/cm². This result stresses again the fact that multi-nucleon emission processes play a significant role

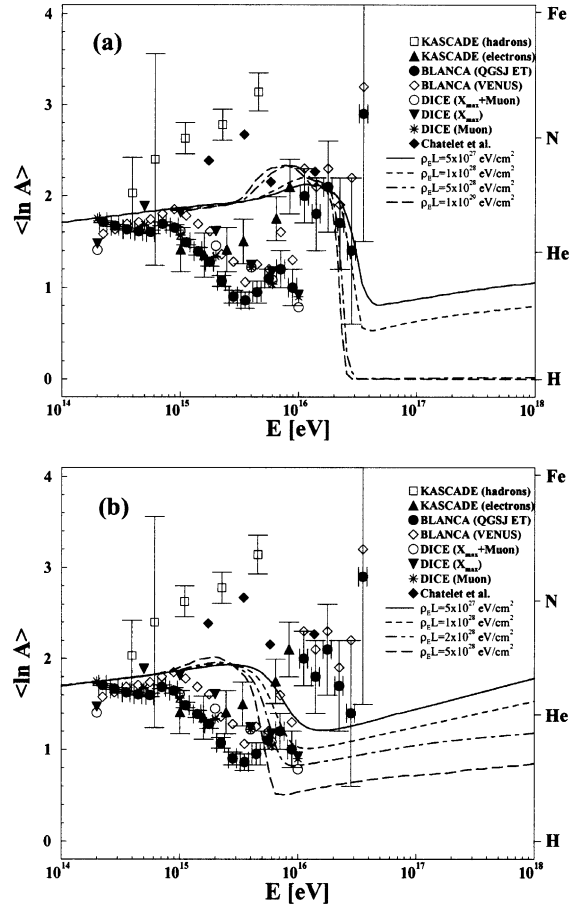


Fig. 4. Plots of mean mass composition $\langle \ln A \rangle$ versus E corresponding to the Planckian photon spectra with (a) $k_B T = 1.8$ eV and (b) $k_B T = 10$ eV. For comparison, the CR mass composition measured by different experiments is also shown. Notice that several data sets exhibit the same qualitative behavior, since they correspond to the same experimental data and differ only in the hadronic models used to interpret the observations.

in CR propagation within this scenario. We also see that, although the solutions with lower column density ($\rho_E L \leq (\rho_E L)_{\max}$) are in reasonable agreement with some data sets within error bars, the composition beyond the knee seems to be somewhat suppressed. Nevertheless, we should bear in mind that we have employed just one source with a fixed column density, and that in a realistic situation there would be many sources, each with a different environment of radiation. Avoiding any

detailed calculations, we may think about the contribution of one additional source. If it happens to have a column density $\rho_E L$ somewhat smaller, it would lead to a flatter flux; then, whereas the first source would be responsible for the occurrence of the knee, the second one would keep a non-vanishing nuclear component beyond the knee. Furthermore, in the computations we assumed that all nuclei traverse the same distance L , thus disregarding the fact that the magnetic field present in the source region should retain more effectively the heavier nuclei at any given CR energy. Taking this into account (for instance, assuming that a nucleus traverses a distance in the photon background proportional to its charge) we would find that, while the disintegration would be very effective for heavy nuclei (and would be, in fact, responsible for forming the knee), it would be less effective for the lighter nuclei (since the column density for them would be much smaller) and they could then survive beyond the knee. Hence, we see that this scenario can easily account for a change in the composition towards lighter nuclei above the knee, which is apparent in some of the existing data. On the contrary, rigidity dependent scenarios predict the opposite trend, and may have difficulties to account for some measurements. It has indeed been pointed out that they would require the introduction of an ad hoc additional light component above the knee to reproduce the observations [24].

The results for the power-law photon spectrum (with a lower cutoff energy in the infrared range) should not differ much from the ones obtained for the Planckian spectrum, and this behavior was in fact corroborated. Let us fix the spectral index and the upper cutoff energy at the values found in Ref. [15] to correspond to the best fit to experimental data, namely $\alpha = 1.3$ and $\epsilon_M = 20$ eV. Fig. 5 shows ϕ_{total} versus E for power-law distributions with column density $\rho_E L = 10^{28}$ eV/cm² and several values for the lower cutoff energy ($\epsilon_m = 0.01, 0.1, 0.8$ eV). It is clear that the overall spectrum displays a steepening which has the right qualitative features to explain the knee although it does not provide an accurate fit to neither AKENO nor CASABLANCA data individually. Turning now to discuss the dependence of the results on the

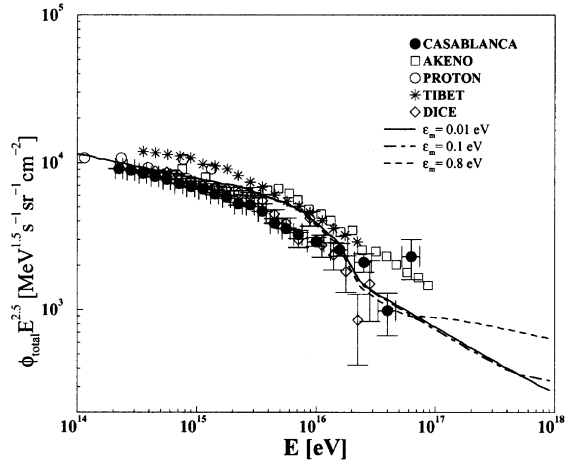


Fig. 5. Plots showing the total CR flux versus the energy per particle for power-law spectra with $\alpha = 1.3$, $\epsilon_M = 20$ eV, $\rho_E L = 10^{28}$ eV/cm² and several values for the lower cutoff energy ($\epsilon_m = 0.01, 0.1, 0.8$ eV). For comparison, experimental data measured by different experiments are also shown.

lower cutoff energy, it is found that, apart from minor differences arising from spectrum normalization, lower values of ϵ_m yield lower fluxes in the high energy end. The explanation for this is that from the expression for the photodisintegration emission rate (Eq. (6)) it turns out that the role of ϵ_m is that of determining a value of the particle energy E , such that below E the giant dipole resonance dominates, while above E the flat high energy regime of σ prevails. As we lower ϵ_m , more photons with lower energy become available so that the flux suppression imposed by the giant resonance extends to higher particle energies.

This analysis provides a suitable explanation of the results shown in Fig. 5, as well as those of Fig. 2(b). The results previously obtained for power-law spectra [15] show an extremely abrupt flux suppression that seems to be in contrast with experimental data (CR energy spectra and, especially, CR mass composition) above $E = 10^{17}$ eV, and can be attributed to the negligible value for ϵ_m adopted there. The use of a non-negligible cutoff prevents the flux suppression and provides a better fit to observations in this energy region. It should also be remarked that a solution yielding an excessive CR flux at energies above 10^{17} eV is by no means troublesome since it is clear that at these

high energies efficient leakage of CRs from the Galaxy cannot be disregarded, and that a less effective acceleration could also be playing a role. These mechanisms would clearly contribute to lower the CR flux, so that it does not seem reasonable to expect a reliable explanation of the CR spectrum by photodisintegration processes alone beyond $E \sim 10^{17}$ eV.

Fig. 6 shows plots of the proton and neutron components ($\phi^p + \phi^{sp}$ and $\phi^n + \phi^{*n}$, respectively) as well as the total CR flux ϕ^{total} for the Planckian photon spectrum with $k_B T = 1.8$ eV and $\rho_E L = 2 \times 10^{29}$ eV/cm². In order to estimate the possible effects of the neutron mechanism, we also plotted the total CR flux that is obtained ignoring the escape of neutrons, labelled ϕ_{total}^* . From the figure it is seen that the neutron mechanism raises the CR total flux by $\sim 100\%$ for $E \geq 5 \times 10^{16}$ eV. Had we adopted a column density $\rho_E L = 10^{29}$ eV/cm², the corresponding flux increment would be only $\sim 10\%$. Nevertheless, it should be noted that these results correspond to very high column densities (actually inconsistent with the observed CR composition), and that effects arising from the neutron escape mechanism turn out to be negligible for lower values of the column density $\rho_E L$, for which also the photopion losses of protons can be altogether ignored.

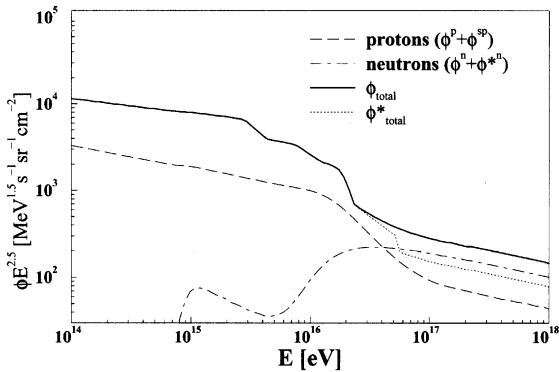


Fig. 6. Plots showing the proton flux ($\phi^p + \phi^{sp}$), the neutron flux ($\phi^n + \phi^{*n}$) and the total CR flux (ϕ^{total}) versus the energy per particle for a thermal photon spectrum with $k_B T = 1.8$ eV and $\rho_E L = 2 \times 10^{29}$ eV/cm². Also shown is the total CR flux (ϕ_{total}^*) that is obtained ignoring the escape of neutrons.

To see what are the implications for the sources in order to produce the required photon densities for this scenario to work, let us make the following simple estimate. An isotropic source with photon energy density $\rho_{E,0}$ at its surface (with radius r_0) will have around it a photon energy density given by $\rho_E(r) = \rho_{E,0}(r_0/r)^2$. This means that the photon column density measured radially from the surface is $\langle \rho_E r \rangle \simeq \rho_{E,0} r_0$. On the other hand, if the CRs are magnetically confined near the source, the path they travel across the dense photon field will be larger, i.e., $\langle \rho_E L \rangle \equiv \lambda \langle \rho_E r \rangle$, with the parameter $\lambda \gg 1$. If to obtain an estimate of the maximum achievable densities we assume that the source radiates with a luminosity $L = 4\pi r_0^2 \rho_{E,0} c$ smaller than the Eddington limit, i.e., $L < L_{\text{Edd}} = 4\pi G M m_p c / \sigma_T \simeq 1.3 \times 10^{38} M/M_\odot$ erg/s, with M the mass of the source, m_p the proton mass and σ_T Thomson's cross-section, we find that the radius of the source needs to satisfy

$$r_0 < \frac{G M m_p}{\sigma_T \rho_{E,0} r_0} \simeq \left(\frac{\lambda}{10} \right) \left(\frac{10^{28} \text{ eV/cm}^2}{\langle \rho_E L \rangle} \right) \left(\frac{M}{M_\odot} \right) 2 \times 10^9 \text{ m}, \quad (29)$$

and hence the photon source must be compact (like a pulsar) and in general smaller than the typical size of a white dwarf. It has to be mentioned however that in non-steady-state situations like in supernova explosions, luminosities larger than L_{Edd} are possible. Another consideration is that if the source has a black-body spectrum with temperature T and it is radiating at the Eddington limit, the relation $4\pi r_0^2 \sigma T^4 = L_{\text{Edd}}$ implies that $T \simeq 1.8 \times 10^3 \text{ eV} (M/M_\odot)^{1/4} (10 \text{ km}/r_0)^{1/2}$, and hence for the photon energies to be in the optical/UV region, radii somewhat larger than those of neutron stars would be favored (or the luminosity should be smaller than L_{Edd}).

As a summary, we have considered in more detail the photodisintegration of CR nuclei in the scenario in which their interaction with optical and soft UV photons around the source is responsible for the steepening of the spectrum and for the change in composition above the knee. The more accurate treatment of the nuclear processes, in

particular the inclusion of the multinucleon emission, implies that the required photon column densities are lower (by an order of magnitude) than previously obtained. The main difference between this scenario and other (rigidity dependent) explanations proposed for the knee is the resulting CR composition, which here becomes lighter above $\sim 10^{16}$ eV.

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References

- [1] J.W. Fowler et al., astro-ph/0003190.
- [2] S.P. Swordy, D.B. Kieda, *Astropart. Phys.* 13 (2000) 137.
- [3] M. Amenomori et al., Mt Fuji Emulsion Chamber Collaboration, in: Proceedings of 19th International Cosmic Ray Conference, vol. 2, La Jolla, 1985, p. 206.
- [4] M. Aglietta et al., EAS-TOP Collaboration, *Nucl. Phys. Proc. Suppl.* 85 (2000) 318.
- [5] K.H. Kampert et al., Cascade Collaboration, OG. 1.2.11 Proceedings of 26th ICRC, Salt Lake City, 1999.
- [6] K.H. Kampert et al., astro-ph/0102266.
- [7] F. Arqueros et al., Hegra Collaboration, *Astron. Astrophys.* 359 (2000) 682.
- [8] C.E. Fichtel, J. Linsley, *ApJ* 300 (1986) 474.
- [9] J.R. Jokipii, G.E. Morfill, *ApJ* 312 (1986) 170.
- [10] K. Kobayakawa et al., astro-ph/0008209.
- [11] S.I. Syrovatsky, *Comm. Astrophys. Space Phys.* 3 (1971) 155.
- [12] J. Wdowczyk, A.W. Wolfendale, *J. Phys. G* 10 (1984) 1453.
- [13] V.S. Ptuskin et al., *Astron. Astrophys.* 268 (1993) 726.
- [14] A.M. Hillas, in: Proceedings of 16th International Cosmic Ray Conference, vol. 8, Kyoto, 1979, p. 7.
- [15] S. Karakula, W. Tkaczyk, *Astropart. Phys.* 1 (1993) 229.
- [16] F.W. Stecker, M.H. Salamon, *ApJ* 512 (1999) 521.
- [17] V.S. Berezinsky, O.F. Prilutsky, *Sov. Astron. Lett.* 3 (3) (1977) 141.
- [18] V.S. Berezinsky, Proceedings of the 15th ICRC, vol. 10, Plovdiv, 1977, p. 84.
- [19] B. Wiebel-Sooth et al., *Astron. Astrophys.* 330 (1998) 389.
- [20] J.L. Puget et al., *Astrophys. J* 205 (1976) 638.
- [21] D.E. Groom et al., *Rev. Part. Phys., Eur. Phys. Jour. C* 15 (2000).
- [22] M. Amenomori et al., *ApJ* 461 (1996) 408.
- [23] E. Chatelet et al., in: Proceedings of 22nd International Cosmic Ray Conference, vol. 2, Dublin, 1991, p. 45.
- [24] S.V. Ter-Antonyan, L.S. Haroyan, hep-ex/000306.