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Large scales anisotropies of extragalactic cosmic rays

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Abstract

We discuss the anisotropies at large angular scales expected for extragalactic cosmic rays. The transition from the 'low' energy galactic cosmic rays to the high energy extragalactic ones probably occurs around EeV energies. Depending on the amplitude of the extragalactic magnetic field and the charge of the particles the propagation of extragalactic cosmic rays from their sources to our galaxy would be in the diffusive regime or the quasi-rectilinear one, with just angular diffusion at higher energies. The anisotropy from a single source of protons is obtained using both numerical simulations and analytical approximations. For a diffusive scenario with a few sources in the local supercluster, we discuss the possible transition between the case in which the anisotropies are dominated by a few sources at energies below few EeV towards the regime in which many sources contribute at higher energies. The effect of a non-isotropic source distribution is also discussed, showing that it can significantly affect the observed dipole.

Keywords:
Ultra high energy cosmic rays, propagation, anisotropies

1. Introduction

The sources of the cosmic rays (CRs) are still unknown, although it is believed that the vast majority of those observed with energies $E < 0.1$ EeV (where EeV $\equiv 10^{18}$ eV) are accelerated in the Galaxy, probably in supernova explosions, while those with energies above few EeV are most likely of extra-galactic origin, probably accelerated in active galaxies or gamma ray bursts. The exact energy at which the transition from galactic to extra-galactic CRs takes place is under debate, with different proposals locating it near the second-knee (a steepening of the spectrum at $\sim$ 0.1 EeV) or near the ankle (a flattening of the spectrum observed at $\sim$ 4 EeV).

Regarding the CR composition, both the maxima in the air shower development as well as the shower maxima fluctuations, determined using the fluorescence technique [1, 2], suggest that at EeV energies CRs are predominantly light nuclei. A proton component at EeV energies can hardly be of galactic origin, since besides the difficulty of devising appropriate galactic sources that could accelerate particles to such high rigidities, the galactic protons would be expected to give rise to strong anisotropies towards the direction of the galactic plane, leading to significant dipolar and quadrupolar components in large scale anisotropy searches, contrary to observations [3]. The Pierre Auger Observatory reported upper bounds on the large scale anisotropies in right ascension of $\sim 2\%$ at EeV energies. The data show a marginally significant indication of a transition from a direction near $RA \simeq 270°$ below 1 EeV, which is consistent with the direction of the galactic center, towards directions near $RA \simeq 100°$ above the ankle energy, with an amplitude increasing to the several percent level near 10 EeV.

It is generally thought that cosmic rays of extragalactic origin at energies below those affected by the Greisen-Zatsepin-Kuzmin (GZK) suppression should be very isotropic because they can reach the observer from cosmological distances and the matter distribution is very homogeneous at large distances. The Earth motion with respect to the rest frame of the matter distribution is expected to produce a small dipolar anisotropy of amplitude $\sim 0.6\%$ [4]. However, if the propagation of CRs in the extragalactic magnetic field is diffusive, as it probably is at EeV energies for light composition and also at much higher energies if CRs are heavy nuclei, larger anisotropies are expected, as we will discuss here following ref. [5].

2. Turbulent magnetic fields and diffusive propagation

Extragalactic magnetic fields are poorly known, they are usually described by a root mean square strength $B = \sqrt{\langle B^2(x)\rangle}$ ranging from nG up to 100 nG and a coherence length with typical values of order $l_c \sim 0.1 - -1$ Mpc [6]. The distribution of the magnetic energy density w on different length scales can be described adopting a power law in Fourier space $w(k) \propto k^{-m}$. For instance, a Kolmogorov spectrum of turbulence corresponds to $m = 5/3$.

For charged particles propagating in a turbulent magnetic field an effective Larmor radius can be introduced as

$$r_L = \frac{E}{ZeB} \simeq 1.1 \frac{E/\mathrm{EeV}}{Z\,B/\mathrm{nG}}\,\mathrm{Mpc}, \quad (1)$$

with Ze the particle charge. A relevant quantity to characterize the particle diffusion is the critical energy E_c, defined such that $r_L(E_c) = l_c$ and hence given by

$$E_c = ZeBl_c \simeq 0.9Z\,\frac{B}{\mathrm{nG}}\,\frac{l_c}{\mathrm{Mpc}}\,\mathrm{EeV}. \quad (2)$$

For energies below the critical energy resonant diffusion takes place, with particles experiencing large deflections induced by their interactions with the B field modes with scales comparable to the Larmor radius. For energies above E_c non-resonant diffusion takes place, in which the deflections after traversing a distance l_c are small, typically of order $\delta \simeq l_c/r_L$.

If diffusion is isotropic the particle flux is given by $\vec{J} = -D\vec{\nabla}n$, with n the particle density and D the diffusion coefficient. A diffusion length can be defined as $l_D \equiv 3D/c$, representing the distance after which the deflection of the particles is $\simeq 1$ rad. For $E \ll E_c$ it is given by $l_D \simeq a_L l_c (E/E_c)^\alpha$, with $\alpha \equiv 2 - m$, while for $E \gg E_c$ it is $l_D \simeq 4l_c(E/E_c)^2$, since in the latter regime one needs to traverse $N \simeq l_D/l_c$ coherent domains to have a total deflection $\delta \simeq 1$ rad, where $\delta \simeq \sqrt{N}(l_c/r_L)$ results from the random angular diffusion of the CR trajectory. From the results of extensive numerical propagation of protons a fit to the diffusion coefficient as a function of the energy was obtained in [5] as

$$D(E) = \frac{c}{3} l_c \left[4\left(\frac{E}{E_c}\right)^2 + a_I \left(\frac{E}{E_c}\right) + a_L \left(\frac{E}{E_c}\right)^{2-m} \right]. \quad (3)$$

For a Kolmogorov spectrum ($m = 5/3$) the coefficients result $a_L \approx 0.23$ and $a_I \approx 0.9$.

Spatial diffusion of the CR particles takes place whenever the distance to the source r_s is much larger than l_D. At sufficiently large energies l_D becomes larger than r_s and the quasi-rectilinear regime takes place, in which the root mean square deflection of the particles arriving from the source is less than 1 rad, and hence only some angular diffusion occurs but not the spatial diffusion. The quasi-rectilinear regime happens for $r_s > l_D$, and hence $E > E_{rect} \equiv E_c\sqrt{r_s/l_c}$ (where we assumed that $E_{rect} > E_c$ so that $D \propto E^2$, which is indeed the case if $r_s \gg l_c$).

The diffusion of extragalactic CRs can modify the spectrum of the particles reaching the Earth, specially at low energies ($E/Z <$ EeV) due to a magnetic horizon effect [7, 8, 9, 10], and can also be crucial to determine the expected anisotropies, as we discuss below.

In the diffusion regime, the density n of ultra-relativistic particles propagating from a source located at $\vec{x}_s$ in an expanding universe obeys the equation

$$\frac{\partial n}{\partial t} + 3H(t)n - b(E,t)\frac{\partial n}{\partial E} - n\frac{\partial b}{\partial E} - \frac{D(E,t)}{a^2(t)}\nabla^2 n = \frac{Q_s(E,t)}{a^3(t)}\delta^3(\vec{x} - \vec{x}_s), \quad (4)$$

where $\vec{x}$ denotes the comoving coordinates, $a(t)$ is the scale factor of the expanding universe, $H(t) \equiv \dot{a}/a$ is the Hubble constant and $D(E,t)$ is the diffusion coefficient. The source function $Q_s(E,t)$ gives the number of particles produced per unit energy and time. The energy losses of the particles are described by

$$\frac{\mathrm{d}E}{\mathrm{d}t} = -b(E,t), \qquad b(E,t) = H(t)E + b_{int}(E). \quad (5)$$

This includes the energy redshift due to the expansion of the universe and energy losses due to the interaction with radiation backgrounds, that in the case of protons include pair production and photo-pion production due to interactions with the CMB background. The general solution was obtained by Berezinsky and Gazizov [11],

$$n(E) = \int_0^{z_{max}} \mathrm{d}z \left|\frac{\mathrm{d}t}{\mathrm{d}z}\right| Q_s(E_g, z) \frac{\exp\left[-r_s^2/4\lambda^2\right]}{(4\pi\lambda^2)^{3/2}} \frac{\mathrm{d}E_g}{\mathrm{d}E}, \quad (6)$$

where z_{max} is the maximum source redshift (note that redshift has the meaning of time rather than distance) and $E_g(E, z)$ is the original energy at redshift z of a particle having energy E at present ($z = 0$). The source function Q_s will be assumed for definiteness to correspond to a power law spectra, $Q_s(E_g) \propto E_g^{-\gamma}$, up to a maximum energy E_{max}. The Syrovatsky variable is given by

$$\lambda^2(E, z) = \int_0^z \mathrm{d}z \left|\frac{\mathrm{d}t}{\mathrm{d}z}\right| (1+z)^2 D(E_g, z), \quad (7)$$

with $\lambda(E, z)$ having the meaning of the typical distance diffused by CRs from the site of their production with energy $E_g(E, z)$ at redshift z until they are degraded down to energy E at the present time. In the expanding universe

$$\left|\frac{\mathrm{d}t}{\mathrm{d}z}\right| = \frac{1}{H_0(1+z)\sqrt{(1+z)^3\Omega_m + \Omega_\Lambda}}, \quad (8)$$

where we consider $H_0 \simeq 70$ km/s/Mpc for the Hubble constant, $\Omega_m \simeq 0.3$ for the matter content and $\Omega_\Lambda \simeq 0.7$ for the cosmological constant contribution.

The amplitude of the dipolar component of the arrival direction distribution can be obtained from the particle density in eq. (6), valid in the diffusive regime, as

$$\vec{\Delta} = \frac{3\vec{J}}{n} = 3D\frac{\vec{\nabla}n}{n} = \Delta\, \hat{r}_s, \quad (9)$$

where $\hat{r}_s$ indicates that the dipole maximum points in the direction of the source.

For large energies we enter the quasi-rectilinear regime, with the arrival directions appearing increasingly clustered around the source location and the diffusion approximation ceases to be valid. In particular, in the limit of small deflections the dispersion of the arrival directions with respect to the source direction is given (in the static case and without energy losses) by [12, 13]

$$\langle\theta^2\rangle = \frac{(Ze)^2 B^2 l_c r_s}{6E^2} = \frac{r_s}{6l_c}\left(\frac{E_c}{E}\right)^2. \quad (10)$$

In the general case the dipolar component of the anisotropy can be computed as follows. The distribution of the arrival directions $\hat{u}$ of particles from a source at $\vec{r}_s \equiv r_s\hat{r}_s$ only depends on the angle between $\hat{u}$ and $\vec{r}_s$, $\theta = \mathrm{acos}(\hat{u}\cdot\hat{r}_s)$, and can be expanded in Legendre polynomials as

$$\Phi(\hat{u}) = f(\cos\theta) = \Phi_0 + \Phi_1\hat{u}\cdot\hat{r}_s + \ldots. \quad (11)$$

The dipole amplitude is then given by

$$\Delta = \frac{\Phi_1}{\Phi_0} = 3\langle\cos\theta\rangle \quad (12)$$

In the next section we will compute the dipolar anisotropy in the full range of energies and distances to the source, covering the transition from spatial diffusion to the quasi-rectilinear regime using numerical simulations of particles propagating in a turbulent magnetic field. By following many particles of a given energy at $z = 0$ propagating back in time in a turbulent magnetic field, $\langle\cos\theta\rangle$ can be computed as the mean cosine of the angle between the original direction of the CR velocity and the vector describing its position when the particles pass at a comoving distance r_s from the original point. In this way the dipolar anisotropy can be numerically obtained for all arrival energies, and we will be able to match the results from the diffusive and quasi-rectilinear regimes.

3. Simulations of charged particle propagation in a turbulent magnetic field

The evolution of the direction of propagation $\hat{n}$ of particles with charge Ze in the turbulent field is followed by integrating the Lorentz equation

$$\frac{\mathrm{d}\hat{n}}{\mathrm{d}t} = \frac{Zec}{E(t)}\hat{n}\times\vec{B}(\vec{x}, t). \quad (13)$$

We will only consider the case of protons in the following, thus $Z = 1$, although if energy losses can be ignored all results also apply to the case of nuclei by replacing E by E/Z (the inclusion of energy losses in the case of nuclei is complicated by the fact that photo-disintegration processes change the nuclear masses and lead to the production of secondary nucleons). The presence of the magnetic field does not change the magnitude of the velocity (nor the particle energy), it only modifies the propagation direction. The dependence with time appearing in eq. (13) arises due to the redshift in the expanding universe and from energy losses due to the interaction of the protons with the CMB radiation. We also included in eq. (13) a possible evolution of $\vec{B}$ with time.

In the non-resonant regime when the deflection in a distance equal to the coherence length l_c is small, what includes both the quasi-rectilinear propagation regime and the regime of spatial diffusion if the propagation distance is larger than the diffusion length, the propagation can be simulated by numerically integrating a stochastic differential equation as proposed in ref. [12]. The magnetic scatterings of the protons lead to angular diffusion of the propagation direction $\hat{n}$ with an angular diffusion coefficient given by

$$\mathcal{D}_{ij} \equiv \frac{\langle\Delta n_i \Delta n_j\rangle}{2c\Delta t} = \frac{l_c}{8}\left(\frac{ZeB}{E}\right)^2 P_{ij} \equiv \mathcal{D}_0 P_{ij}, \quad (14)$$

where $P_{ij} \equiv (\delta_{ij} - n_i n_j)$ is the tensor projecting to the plane orthogonal to $\hat{n} \equiv (n_1, n_2, n_3)$.

In this framework it is not difficult to take into account the effects of the expansion of the universe and the proton energy losses due to interactions with the CMB radiation [5]. We can include the possible variation of the magnetic field amplitude with the redshift, that we can take as $B(z) = B(0)(1+z)^{2-\mu}$, where the factor $(1+z)^2$ arises from the flux conservation as the universe expands and the index μ was introduced in ref. [11] to account for magneto-hydrodynamic effects, and was taken there as $\mu = 1$, as we will adopt in the simulations. The critical energy dependence with z is $E_c(z) = E_c(0)(1+z)^{1-\mu}$. The coherence length will typically scale as $l_c(z) = l_c(0)/(1+z)$.

We follow protons from the observation time at $z = 0$ back in time with steps adapted such that the change in the comoving coordinates $|\Delta\vec{r}|$ is equal to $l_c(0)$ in each step, $c|\Delta t| = |\Delta\vec{r}|/(1+z)$. This corresponds to a step in redshift $|\Delta z| = l_c(0)\sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}H_0/c$. Due to energy losses from redshift and from the interactions with the CMB particles arriving to $z = 0$ with energy $E(0)$ had a larger energy $E_g(z)$ at each previous step. We take into account these effects by integrating eq. (5) with b_{int} parametrized using an analytic fit to the attenuation length (see [5]). Then, in each step the particle moves a comoving distance $l_c(0)$ and the propagation direction suffers a stochastic change given by

$$(\Delta\hat{n})_i = \sqrt{2\mathcal{D}_0(z)l_c(0)/(1+z)}P_{ij}\xi_j, \quad (15)$$

where

$$\mathcal{D}_0(z) = \frac{1}{8l_c(0)}\left(\frac{E_c(0)}{E_g(z)}(1+z)^{3-2\mu}\right)^2. \quad (16)$$

The new propagation direction is obtained from

$$\hat{n}(z+\Delta z) = \sqrt{1-|\Delta\hat{n}|^2}\,\hat{n}(z) + \Delta\hat{n} \quad (17)$$

and the new comoving coordinate from

$$\vec{x}(z+\Delta z) = \vec{x}(z) + l_c(0)\hat{n}(z). \quad (18)$$

In this way the proton trajectories can be followed back in time. In order to compute the expected dipole anisotropy of particles from a source at comoving distance r_s, for each arrival energy $E(0)$ we backtrack the trajectories of a large number of particles and compute the mean cosine of the angle between the initial direction and the position when the particles pass at a comoving distance r_s from the original point. When taking the mean we have to include a weight factor for each particle equal to $[E_g(z)/E(0)]^{-\gamma}\mathrm{d}E_g/\mathrm{d}E$, where the first factor takes into account that for a source emitting protons with a spectrum $E^{-\gamma}$ there will be less particles with the higher energy $E_g(z)$ required to reach the observer with a given $E(0)$, and the second factor takes into account the change in the energy bin width from the emission to the observation.

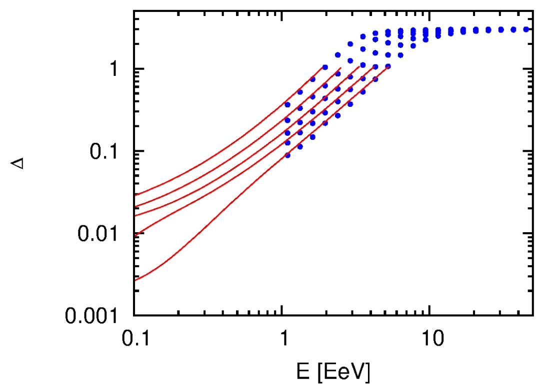

Figure 1: Dipole amplitude for one source located at a comoving distance of 25, 50, 100, 200 and 400 Mpc (from top to bottom) for a coherence length of the magnetic field of $l_c = 1$ Mpc and an amplitude $B = 1$ nG.

Figure 1 shows the dipole anisotropy as a function of the energy arising from one single source located at a comoving distance of 25, 50, 100, 200 or 400 Mpc (from top to bottom). Blue dots show the results obtained from the integration of the stochastic differential equation while solid red lines show the results from the solution to the diffusion equation (using eqs. (9) and (6)). A very good agreement is seen in the overlapping region. A spectral index $\gamma = 2$ and a maximum energy $E_{max} = 10^{21}$ eV are considered in all the examples. The magnetic field coherence length was taken as $l_c = 1$ Mpc and B= 1 nG, leading to $E_c = 0.9$ EeV.

4. Large scale anisotropy from many sources

In the previous section the dipolar anisotropy produced by an individual source in the presence of a turbulent magnetic field was computed as a function of the source distance and of the CR energy. In a realistic situation the total cosmic ray flux will probably originate from a set of several (or many) sources. The total dipolar component of the flux will mainly depend on the location and intensities of the nearest sources and on whether there is an inhomogeneous distribution of the sources at large scales. If there are several sources contributing to the flux, the dipolar anisotropy can be obtained from the superposition of the individual source

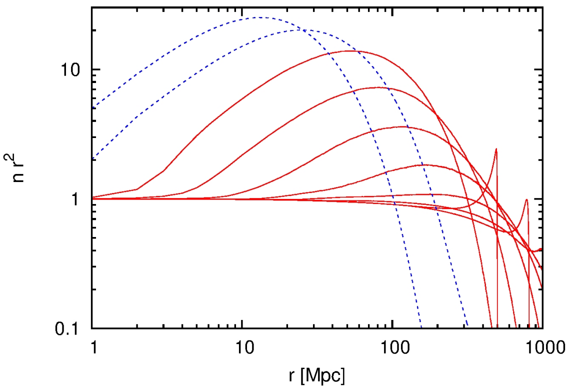

Figure 2: Density of cosmic rays $n(r)$ times r^2 as a function of the distance from the source to the observer for $B = 1$ nG and for different observed energies, $E = 0.9, 1.5, 2.6, 4.4, 7.4, 12.5, 21.2$ and 36 EeV (red solid lines from top-left to bottom-right) and $E = 0.1$ and 0.3 EeV (blue dashed lines).

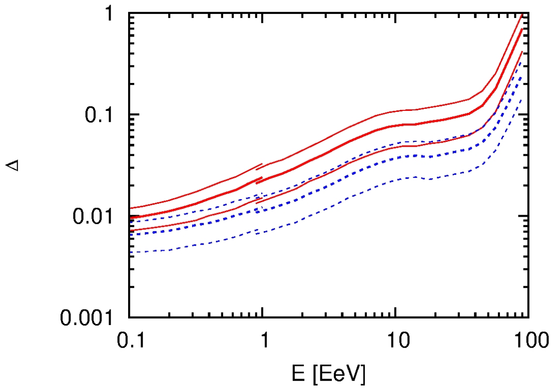

Figure 3: Mean and dispersion of the total dipole amplitude as a function of the energy for a turbulent magnetic field of $B = 1$ nG and a density of sources $\rho = 10^{-5}$ Mpc^{-3} (red solid lines) and $\rho = 10^{-4}$ Mpc^{-3} (blue dashed lines).

dipoles through

$$\vec{\Delta}(E) = \sum_{i=1}^{N} \frac{n_i}{n_t}(E)\vec{\Delta}_i(E), \qquad (19)$$

where N is the number of sources giving a non-negligible contribution to the flux at energy E, n_i/n_t measures the fraction of the flux coming from the i-th source and $\vec{\Delta}_i(E)$ is the dipole anisotropy of the flux from source i computed in the previous section and shown in Figure 1.

In order to estimate the relative contribution to the anisotropy of the different sources as a function of their distance to the observer we will make the simplifying assumption that the sources are steady and have equal intrinsic intensities, so that for each energy the relative contribution to the flux from different sources will only depend on the distance to the source r_i. The product $n(r)r^2$ is shown in Figure 2 for different energies (arbitrary normalization).

At low energies, where particles are in the regime of spatial diffusion, we can see a large enhancement of the flux with respect to the typical $n \propto r^{-2}$ behavior characteristic of rectilinear propagation, followed by a drop at large distances corresponding to the magnetic horizon effect. For large energies the diffusion enhancement disappears as particles travel more straight and they can also arrive from larger distances. For the largest energies the maximum distance from which sources contribute is limited due to energy losses (GZK horizon), and a bump in the flux appears at large distances due to the pile-up caused by the effect of the photo-pion production threshold. We note that even if we assumed the magnetic field turbulence to be uniform everywhere, which may be a crude approximation at very large distances due to the presence of voids and filamentary structures in the matter distribution, at the energies for which distances much larger than $\sim$ 100 Mpc are relevant, the overall contribution from far away sources is not expected to be strongly affected by diffusion effects.

The fact that the sources are distributed in different sky directions means that the vector sum in eq. (19) will generally lead to a smaller dipole amplitude when many sources contribute. In the case that only few sources are relevant, the direction of these particular sources will determine the dipolar anisotropy, while if many sources are relevant, the overall large scale distribution of the sources, in particular whether the distribution has a non-vanishing dipole component, can have a significant effect.

In order to quantify the total amplitude of the dipolar anisotropy we performed some simple simulations. Starting with one source at a random direction in the sky, that represents the closest source, we subsequently added new sources with equal intrinsic CR luminosity in random directions and computed the new total dipolar anisotropy using eq. (19). The radial distances from the observer to the sources are taken as the mean expected value for the i-th closest source in an homogeneous distribution, that is given by $\langle r_i \rangle = (3/4\pi\rho)^{1/3}\Gamma(i + 1/3)/(i-1)!$, where ρ is the density of sources. Figure 3 shows the amplitude of the dipole and the dispersion obtained in 1000 simulations for two different values of the source density: $\rho = 10^{-5}$ Mpc^{-3}, for which the closest source is at a mean distance $\langle r_1 \rangle \simeq 25$ Mpc (solid lines), and $\rho = 10^{-4}$ Mpc^{-3}, for which $\langle r_1 \rangle \simeq 11$

Mpc (dashed lines). For $\rho = 10^{-5}$ Mpc^{-3} the dipole amplitude rises from about 1% at $E = 0.1$ EeV to $\sim$ 2% at 1 EeV and $\sim$ 8% at 10 EeV. For $\rho = 10^{-4}$ Mpc^{-3} an anisotropy smaller by a factor of about 2 results, as many more sources contribute to the total flux in this case. At the largest energies, above the GZK cutoff, a steep increase of the anisotropy results as the number of contributing sources decreases. Note that the computations are performed using the stochastic integration results for energies above E_c, while below E_c the diffusion solution from Section 2 is used. This explains the small jumps observed in the plots at $E \simeq E_c$.

The previous results hold for homogeneously distributed sources. In the case that the sources themselves have an inhomogeneous distribution around the observer, in particular when the distribution has a non-vanishing dipole, a further contribution to the anisotropy is expected. Actually, if the local distribution of cosmic ray sources follows the local distribution of matter, a non-vanishing dipole is expected. This dipolar component of the matter distribution is indeed known to be responsible for the Local Group peculiar velocity with respect to the rest frame of the CMB, that gives rise to the observed CMB dipole. The dipolar component of the mass distribution in our neighborhood has been estimated using different catalogs of galaxies as for example the 2 Micron All-Sky Redshift Survey (2MRS), showing that the resulting dipole seemingly converges when sources up to a distance $\sim$ 90 Mpc are included [14]. The effect of the local inhomogeneity of the source distribution in the predicted large scale anisotropies can be included in the simulations by choosing the positions of the sources in our neighborhood from some catalog representing the local distribution of matter. To estimate the effect, we present the results using a volume limited subsample[1] of the 2MRS catalog up to 100 Mpc [15]. We have then selected the position of the required number of sources (according to the density considered) from this subsample of 2MRS galaxies. On the other hand, the locations of sources farther away than 100 Mpc were assumed to be isotropically distributed. We show in Figure 4 the mean dipole amplitude when the inhomogeneous source distribution is considered for a density $\rho = 10^{-5}$ Mpc^{-3} and a turbulent field of $B = 1$ nG. An enhancement of the dipole amplitude of about 70% on average is observed with respect to the isotropic case.

[1]Considering only objects with $d < 100$ Mpc and absolute magnitude in the K band $M_K < -23.4$

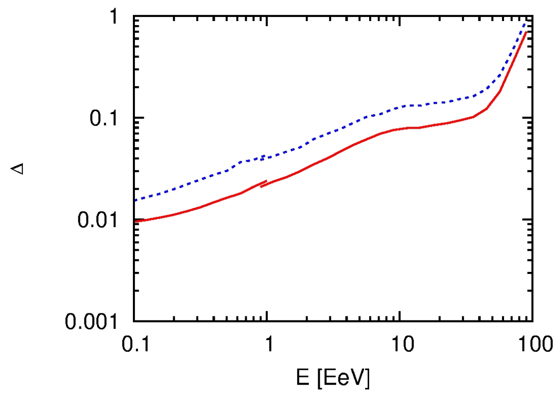

Figure 4: Mean amplitude of the total expected dipole when local sources within 100 Mpc are distributed like galaxies in the 2MRS catalog (blue dashed lines) considering a density $\rho = 10^{-5}$ Mpc^{-3} and a turbulent field with $B = 1$ nG. For reference, the red line shows the expected amplitude for uniformly distributed sources for the same parameters.

5. Summary and discussion

The diffusion of ultra high energy protons in turbulent magnetic fields has been studied, focusing in the computation of the dipolar anisotropies. We obtained the expected large scale anisotropies of extragalactic protons for an extended energy range, matching the results of the high energy regime of angular diffusion (quasi rectilinear propagation) with the low energy regime of spatial diffusion. We illustrated the results for typical values of $E_c \simeq 1$ EeV and $l_c \simeq 1$ Mpc, showing that the dipole amplitude resulting from sources with number densities of 10^{-5} to 10^{-4} Mpc^{-3} are at the level of (0.5–1)% at 0.1 EeV energies, increasing to (1–2)% at 1 EeV and up to (3–10)% at 10 EeV. When the anisotropy in the local (within 100 Mpc) distribution of sources, modeled following the 2MRS galaxy catalog, is taken into account, an increase in the expected dipole amplitude typically by a factor 1.5 to 2 is predicted. In this case this contribution would point in the approximate direction of the motion of the Local Group with respect to the CMB rest frame[2], since it is just the anisotropy in the galaxy distribution that is ultimately responsible for both the proper motion of the Local Group and for the anisotropy in the CR source distribution. These anisotropies are significantly larger than the ones that would result from the Compton-Getting effect [4] if CRs were isotropic in the rest frame of the CMB, which would be at the level of $\sim$ 0.6% almost independently of the energy.

[2]In the rest frame of the CMB the Local Group moves towards the direction $(\alpha, \delta) = (163°, -27°)$ [14].

It should be mentioned that the further deflections of the CRs caused by the galactic magnetic field (mostly by the regular component), not included in this work, would modify the CR dipole amplitude and direction as well as generate higher order multipoles in the arrival directions distribution (see [16]). This effect could reduce to some extent the amplitudes obtained here for energies below few EeV.

The large scale anisotropy of extragalactic cosmic rays obtained are of interest to interpret the recent results obtained by the Pierre Auger Observatory hinting at non-vanishing dipolar amplitudes at energies above $\sim$ 1 EeV [3]. We should notice that the dipolar anisotropies were obtained under the assumption of a proton CR composition (which is consistent with the observations at EeV energies and also compatible with HiRes and Telescope Array measurements at higher energies [2]). If the CR composition were to become heavier above a few EeV, as suggested by the Auger Observatory measurements [1], the anisotropies above the ankle will also depend on the details of the actual source composition.

Acknowledgments

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