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The cosmic-ray air-shower signal in Askaryan radio detectors

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ABSTRACT

We discuss the radio emission from high-energy cosmic-ray induced air showers hitting Earth's surface before the cascade has died out in the atmosphere. The induced emission gives rise to a radio signal which should be detectable in the currently operating Askaryan radio detectors built to search for the GZK neutrino flux in ice. The in-air emission, the in-ice emission, as well as a new component, the coherent transition radiation when the particle bunch crosses the air–ice boundary, are included in the calculations.

1. Introduction

We calculate the radio emission from cosmic-ray-induced air showers as a possible (background) signal for the Askaryan radio-detection experiments currently operating at Antarctica [1–3]. A high-energy neutrino interacting in a medium like (moon)-rock, ice, or air will induce a high-energy particle cascade. In 1962 Askaryan predicted that during the development of such a cascade a net negative charge excess arises mainly due to Compton scattering [4]. This net excess charge by itself will induce a radio signal that can be used to measure the original neutrino. This Askaryan radio emission [4–6] has been confirmed experimentally at SLAC [7], and more recently the Askaryan effect was also confirmed in nature by the radio emission from cosmic-ray induced air showers [8–10].

For high-energy cosmic-ray air showers, along with the Askaryan emission, there is another emission mechanism due to a net transverse current that is induced in the shower front by Earth's magnetic field [11–14]. Recently the radio emission from cosmic-ray air showers has been measured in great detail by the LOFAR collaboration [10,15,16], confirming the predictions from several independent radio emission models [17–20].

Most Askaryan radio detectors [1–3,21–23] search for so-called GZK neutrinos that are expected from the interaction of ultra-high-energy cosmic-ray protons with the cosmic microwave background [24,25]. The expected GZK neutrinos are extremely energetic with energies in the EeV range, while the flux at these energies is expected to fall below one neutrino interaction per cubic kilometer of ice per year. Therefore, to detect these neutrinos an extremely large detection volume, even larger than the cubic kilometer currently covered by the IceCube experiment, is needed. Due to its long attenuation length, the induced radio signal is an excellent means to detect these GZK neutrinos. This has led to the development of several radio-detection experiments [1–6,26–30]. Nevertheless, the highest-energy neutrinos detected so-far are those observed recently by the IceCube collaboration [31] and have energies up to several PeV, just below the energies expected from the GZK neutrino flux.

In this article we calculate the radio emission from cosmic-ray-induced air showers as a possible (background) signal for the Askaryan radio-detection experiments currently operating at Antarctica [1–3]. Besides the emission during the cascade development also transition radiation should be expected when the cosmic ray air shower hits Earth's surface [32,33]. It follows that the induced emission is very hard to distinguish from the direct Askaryan emission from a high-energy neutrino induced cascade in a dense medium such as ice.

2. Radio emission from a particle cascade

We start from the Liénard–Wiechert potentials for a point-like four current from classical electrodynamics and closely follow the macroscopic MGMR [34] and EVA [20] models. Both models were developed to describe the radio emission from cosmic-ray-induced

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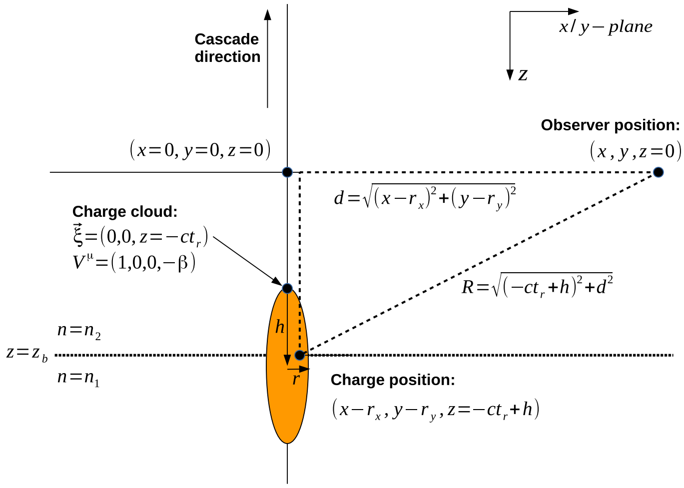

Fig. 1. The geometry used to calculate the radiation emitted from a charge cloud crossing a boundary at $z = z_b$. The observer is positioned at an impact parameter $d = \sqrt{(x - r_x)^2 + (y - r_y)^2}$.

air showers. The Liénard–Wiechert potentials for a point charge, $A^\mu_{PL}(t, \vec{x})$, as seen by an observer positioned at $\vec{x}$ at an observer time t are obtained directly from Maxwell's equations after fixing the Lorenz gauge [35],

$$A^\mu_{PL}(t, \vec{x}) = \frac{1}{4\pi\epsilon_0}\frac{J^\mu}{|\mathcal{D}|}\bigg|_{ret}. \tag{1}$$

The point-like current is defined by $J^\mu = eV^\mu$, where e is the charge, and V^μ is the four-velocity for a particle at $\vec{\xi}(t_r)$ where the retarded emission time is denoted by t_r. The denominator of the vector potential, $\mathcal{D}$, is the retarded four-distance. For an extended current with longitudinal dimension h and lateral dimensions $\vec{r}$, the vector potential has to be convolved with the charge distribution given by the weight function $w(\vec{r}, h)$,

$$A^\mu(t, \vec{x}) = \frac{1}{4\pi\epsilon_0}\int \mathrm{d}h\, \mathrm{d}^2r\, \frac{J^\mu w(\vec{r}, h)}{|\mathcal{D}|}\bigg|_{ret}, \tag{2}$$

where the vector potential has to be evaluated at the retarded emission time t_r. The corresponding geometry is denoted in Fig. 1. We consider an observer positioned at an impact parameter $d = \sqrt{(x - r_x)^2 + (y - r_y)^2}$ perpendicular to the charge track, where r_x, and r_y denote the lateral position of the considered charge within the charge cloud. Defining the element in the plane of the observer perpendicular to the charge-track as $z = 0$, we can define the time at which the front of the charge cloud crosses this plane to be $t = 0$. Using these definitions the position of the charge along the track is now given by $z = -ct_r + h$.

Fixing the geometry, the vector potential can now be evaluated. The retarded emission time is obtained from the light-cone condition with respect to the optical path length L,

$$c(t - t_r) = L, \tag{3}$$

from which the relation between the observer time and the emission time, $t_r(t)$, can be obtained. It should be noted that t_r is a negative quantity. For a medium consisting out of m layers with different index of refraction n_i, the optical path length can be defined by

$$L = \sum_{i=1}^{m} n_i d_i, \tag{4}$$

where the distance d_i, the distance covered by the emission in layer i, is obtained by using a ray-tracing procedure based on Snell's law. Following [36], the retarded distance for a signal traveling through different media is given by,

$$\mathcal{D} = L\frac{\mathrm{d}t}{\mathrm{d}t_r}. \tag{5}$$

In this work the index of refraction is assumed to be independent of frequency within the radio frequency range starting from a few MHz, up to several GHz. In the simplified situation where the signal travels through a medium with constant index of refraction n, the retarded distance can be written in the more common form,

$$\mathcal{D} = nR(1 - n\beta\cos(\theta)), \tag{6}$$

where θ denotes the opening angle between the line of sight from the emission point to the observer and the direction of movement of the emitting charge.

2.1. Cherenkov effects for a single electron

For a single electron moving at a highly relativistic velocity $\vec{\beta} = \vec{v}/c \approx 1$ along the z-axis (by definition), the current is given by $J^\mu = e(1, 0, 0, -\beta)$. The electric field is now obtained directly from the Liénard–Wiechert potentials through,

$$E^i(t, \vec{x}) = -\frac{\mathrm{d}A^0}{\mathrm{d}x^i} - \frac{\mathrm{d}A^i}{\mathrm{d}ct}, \tag{7}$$

where $i = x, y$ gives the polarization of the field in the transverse direction, and x^i denotes the observer position in the transverse plane ($x^1 = x, x^2 = y$). For the moment we will ignore the electric field in the longitudinal direction and, since $A^i \propto J^i = 0$ for $i = 1, 2$ (there is no transverse current), we only have to consider the spatial derivative of the scalar potential. The electric field in the longitudinal direction will in general be small and can easily be calculated following the gauge condition $\vec{k}\cdot\vec{\epsilon} = 0$, where $\vec{k}$ is the momentum vector of the photon and $\vec{\epsilon}$ the polarization. Hence the photon cannot be polarized along its direction of motion. Starting at the zeroth component of the vector potential, the spatial derivative can be evaluated by,

$$\frac{\mathrm{d}A^0}{\mathrm{d}x^i} = \frac{\partial}{\partial x^i}A^0, \tag{8}$$

which corresponds to the radiation from a net charge moving through the medium. For a relativistic electron ($\beta \approx 1$) moving in a medium with a refractive index $n > 1$ this term becomes,

$$E^i_{st}(t, \vec{x}) = -\frac{\partial}{\partial x^i}A^0 = \frac{-e}{4\pi\epsilon_0}\frac{(1 - n^2)x^i}{|\mathcal{D}|^3}, \tag{9}$$

where the label 'st', denotes that the field is due to a highly relativistic non time-varying steady charge. The emission shows a radial polarization direction and vanishes linearly with the distance of the observer to the shower core. This component of the electric field is suppressed by the factor $1 - n^2$, which vanishes in vacuum. In a medium with an index of refraction larger than unity, however, this factor does not vanish and Cherenkov radiation is observed at the point where the retarded distance vanishes, $\mathcal{D} = \sqrt{t^2 + (1 - n^2\beta^2)(x^2 + y^2)} = 0$.

The retarded distance vanishes at the finite Cherenkov angle $\cos(\theta_{CH}) = \frac{1}{n\beta}$ (see Eq. (6)) where the electric field diverges. One intuitive way to understand the Cherenkov effect follows from the more general definition of $\mathcal{D}$ given in Eq. (5). For a vanishing retarded distance, the derivative $\mathrm{d}t/\mathrm{d}t_r$ has to vanish. It follows that the function $t(t_r)$ is flat at this point. Hence at an observer time t, signals emitted at different emission times t_r will be observed at once, leading to a boosted electric field. The vanishing of the retarded distance leads to a divergence in the electric field expressions. These divergences are integrable and therefore disappear for coherent emission by performing an integration over the finite charge and current distributions in the shower front [20].