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Generation of narrow peaks in spectroscopy of charged particles

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In spectroscopy of charged particles, narrow peaks may appear in continuous spectra if magnetic transport of the particles is involved. These artefacts, which so far have escaped the attention of investigators, can develop whenever geometric detection efficiency is less than 100%. As such peaks may be misinterpreted as new physics, their generation is investigated, both analytically and experimentally, for various detector configurations, including those used in searches for the spontaneous decay of the vacuum in heavy-ion collisions.

1. Introduction

New phenomena are often discovered as unexpected peaks in otherwise smooth spectra. Therefore it is of interest to identify sources of artificial peaks that may lead to wrong conclusions. The purpose of this article is to show how easy it is to produce narrow peaks in continuous spectra from seemingly well understood instrumental configurations.

We treat the case, widely met in experimental science and applications, of a point source, placed in a magnetic field that guides the emitted charged particles to a detector, or to some other target, see Fig. 1.

We first summarize recent theoretical and experimental results on unexpected singularities that appear in the so-called magnetic point spread function (PSF), defined for a monoenergetic source. We present some improvements of our theory (Section 2), as well as a measurement of the magnetic PSF in a simple experiment (Section 3). Based on this, we develop the theory of energy spectra in the presence of magnetic transport, and show that singularities in the PSF may lead to prominent structures in an initially smooth continuous spectrum. Predictions based on this work are tested experimentally (Section 4). Finally, we apply our findings to past experiments in positron spectroscopy, which searched for the spontaneous decay of the vacuum in deep inelastic heavy-ion collisions (Section 5).

2. The magnetic point spread function of a monoenergetic source

Recently, the spectroscopy of charged particles was investigated in the context of neutron β-decay [1]. It was shown that, following magnetic transport from a monoenergetic point source to a detector, an infinite number of ring-shaped singularities of radii R_n appear in the source's image on the detector, as shown in Fig. 2. In optics, the image of a point source is called the point spread function of the system. It provides the intensity distribution $f(R)$ in the image plane as a function of the distance R from the optical axis. In our case of magneto-optics, we call this function the magnetic PSF. The singularities in the magnetic PSF were investigated in more detail in [2]. The topic was further developed in [3], while experimental effects of these singularities were first measured in [4]. Although the derivation of the true magnetic PSF and its singularities is rather simple, it seems to have escaped the attention of investigators. A first mention of singularities in the context of a one-turn magnetic spectrometer in Ref. [5] apparently had fallen into oblivion.

The gyration radius of a particle, with charge e and mass m in a magnetic field of amplitude B, is

$$r = r_0 \sin\theta, \text{ with } r_0 = p/eB, \tag{1}$$

with relativistic momentum $p = c^{-1}(E^2 + 2mc^2E)^{1/2}$, and polar angle θ with respect to the field direction. Following magnetic transport over a distance z_0 in a uniform field, the particle reaches the detector after a number of n' helical orbits of gyration, with $n' = \alpha/2\pi = z_0/d$ (we reserve the undashed n for integer values of n'). Here α is the phase angle of the gyrating particle, and $d = 2\pi r_0 \cos\theta$ is

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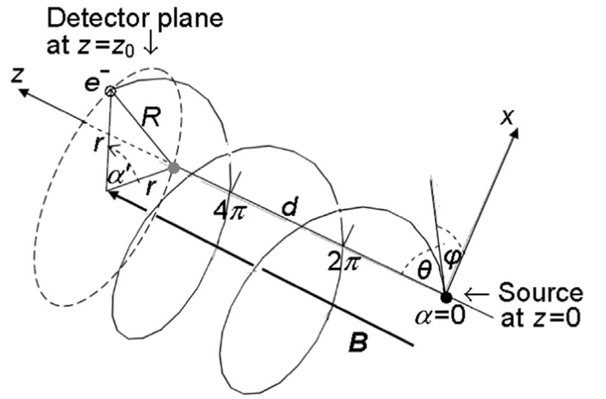

Fig. 1. A point source of charged particles at $\boldsymbol{x}=0$, and a flat detector at $z=z_0$, are both coupled by a uniform magnetic guide field B applied along z. Indicated are the polar and azimuthal emission angles θ and φ, the radius of gyration r, the pitch of the helix d, the phase angle α, with $\alpha'=\alpha$ modulo 2π, and the displacement R of the point of impact on the detector from the central point of impact (the gray dot, reached for emission under $\theta=0$).

the pitch of the helix (Fig. 1). Hence, when the particle hits the detector, its phase angle α on the detector surface (relative to an azimuthal emission angle φ) depends on the emission angle θ at the source as

$$\alpha = z_0/(r_0 \cos\theta). \tag{2}$$

The displacement of the particle from the central point of impact (the gray dot in Fig. 1) depends on α as

$$R(\alpha) = 2r|\sin 1/2\alpha| = 2r_0(1-\alpha_0^2/\alpha^2)^{1/2}|\sin 1/2\alpha|, \tag{3}$$

given by simple geometry, cf. Fig. 1. The angle $\alpha_0 = z_0/r_0$ is the smallest occurring phase angle, for emission in the limit $\theta \to 0$. The corresponding smallest number of helical orbits is $n_0=\alpha_0/2\pi$, or

$$n_0 = eBz_0/2\pi p, \tag{4}$$

which is the only parameter of the theory. $R(\alpha)$ from Eq. (3) is plotted in Fig. 3a) for the same n_0 as used in Fig. 2.

The conventional approach to the problem, summarized in [4] (for references see [2]), leads to the smooth hyperbolic PSF $f_{conv}(R)=1/(4\pi R r_0)$, shown in Fig. 2a). In the conventional approach, it is assumed that all phase angles $\alpha'=\alpha$ modulo 2π at the position of the detector surface, Fig. 1, occur with equal probability, disregarding the functional relation Eq. (2) of α and θ. However, a particle emitted with polar angle θ reaches one single point on the detector, whose corresponding phase angle α (relative to the starting angle φ) is fixed by Eq. (2). Different values of α on the detector correspond to different angles θ at the source with a different gyration radii $r = r_0 \sin\theta$. Hence, with increasing θ, and thereby increasing α, the impact positions on the detector move along a spiral. In our figures, all displacements R are shown in units of gyration radius r_0, Eq. (1).

This spiral is seen in Fig. 3b), where the displacement R is shown not as $R(\alpha)$ in Cartesian coordinates as in Fig. 3a), but as $R(\beta)$ in polar coordinates, with $\beta=\alpha/2+\varphi$. The values of R_n shown correspond to the maximum values of R reached when α increases steadily, see Eqs. (6) and (7). When these impact points are averaged over azimuthal emission angles φ, one obtains the rotationally symmetric true PSF $f(R)$. This is the method applied numerically in Ref. [3]. Our analytical method is as follows.

The magnetic PSF is best written as

$$f(R) = \frac{1}{2\pi R}\left|\frac{dP}{d\cos\theta}\frac{d\cos\theta}{d\alpha}\frac{d\alpha}{dR}\right|, \tag{5}$$

with the angular distribution $dP/d\cos\theta$ at the source, where dP is the emission probability into the solid angle element $d\cos\theta$. With $d\cos\theta/d\alpha$ from Eq. (2), and with $R(\alpha)$ and $d\alpha/dR$ from Eq. (3), the PSF is easily calculated as a function of α, see Eq. (13) of [2]. To express the PSF as a function of R, however, one needs to invert $R(\alpha)$, Eq. (3) and Fig. 2a), which is not possible algebraically. This inversion can, in principle, be done numerically for each value of R, but this is quite time consuming, in particular if, as in the following sections, we have to integrate the singularity-ridden PSF over R for many different values of energy E.

We therefore use the approximation described in [1,2], in which $R(\alpha)$, between two of its zeros (roots), is piecewise approximated by invertible cosine functions passing through the maxima R_n of the true $R(\alpha)$ expressed in Eq. (3), which is shown by the dashed curves in Fig. 3a). The lowest cycle (starting at $n=n_0$, for particle emission in the limit $\theta \to 0$), is numbered n_f = floor(n_0), with floor indicating the next integer below n_0. There R has its maximum at

$$R_{n_f} = \text{Max}\,[R(\alpha),\, 2\pi n_0 \le \alpha \le 2\pi(n_f+1)], \tag{6}$$

while for the higher cycles with $n > n_f$,

$$R_n = \text{Max}\,[R(\alpha),\, 2\pi n \le \alpha \le 2\pi(n+1)]. \tag{7}$$

In addition, we introduce two improvements: First, we adapt

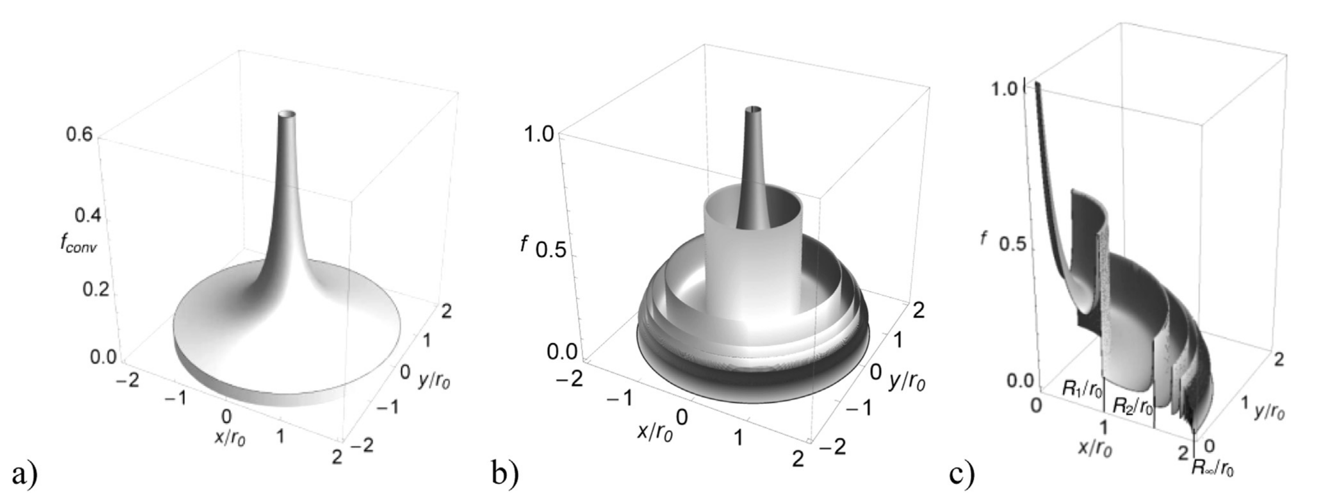

Fig. 2. a) The conventional magnetic PSF on the detector's x-y surface has a smooth hyperbolic shape $f_{conv}(R) \propto 1/R$, with $R=(x^2+y^2)^{1/2}$. **b)** The true magnetic PSF $f(R)$ has an infinite number of singularities. **c)** $f(R)$ in the first quadrant, with radii R_n of the singularities. The PSF is calculated for instrument parameter $n_0=2.4$, Eq. (4). This value for n_0 is obtained, for example, with 1 MeV electrons gyrating through a distance $z_0=0.2$ m in a 0.36 T field, i.e., gyration radius $r_0=$ 1.3 cm.