Detection of X-Ray Transition Radiation of 31-GeV Electrons

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An improved method of particle detection is described, utilizing both the x-ray transition radiation and the ionization of a gas by the particle.

High-energy particle detectors based on x-ray transition radiation (XTR), studied up to the present time, possess small efficiency in particle registration and small relative aperture.[1-5] High registration efficiency was achieved by the streamer-chamber technique,[6] but with this technique it is impossible to form the trigger pulse. We describe here a particle detector which takes advantage of the registration of energy released in the bulk of a gas by the absorption of XTR quanta as well as by the ionization of the gas by a high-energy primary particle. It has been observed from data on the energy release that high registration efficiency of ultrahigh-energy particles may be achieved using the XTR technique, the relative aperture being rather high. This method, first proposed in 1961,[7] was experimentally investigated only recently.[8] It consists in the registration, by a single detector of x-ray quanta, of both the charged particle and the transition radiation generated by the particle in a laminar medium. The yield of the detector will be proportional to the total release of energy due to the absorption of XTR and the ionization losses of the particle in the same detector. While the intensity of XTR strongly depends on $\gamma = E/mc^2$,[9] the ionization losses of ultrarelativistic particles are practically independent of γ; the yield of the detector for small values of γ is defined, consequently, by ionization losses, and for large γ basically by the absorption of XTR.

The registration of radiation and particles in the high-energy particle detector under examination was by means of a gaseous xenon scintillator enclosed in an aluminum container with round, 100-μm-thick Mylar windows, 6 cm in diameter, through which the radiation and registered particles passed (Fig. 1). The photomultiplier, with its photocathode right in the gaseous medium, was inserted from the lateral side of the container, the thickness of the scintillator being 4 cm of xenon at a pressure of 1.6 atm.

As the wavelengths of the light emitted in a gaseous scintillator are in the far ultraviolet, a coating of the inner side of the container and the photocathode by a spectrum transformer was provided to match the radiation spectrum and the spectral characteristic of the photomultiplier. Special attention was given to purifying the gas, since impurities cause a sharp reduction of scintillation intensity. To that end, the scintillation gas was continuously purged in 600°C hot calcium chips by means of natural circulation.

The high-energy particle-detector assembly, consisting of a laminar medium followed by the xenon gaseous scintillator, was exposed to 31-GeV electrons at the Serpukhov proton accelerator.[10] The electron beam was separated by an array of scintillation counters. The laminar medium was composed of 1000 10-μm-thick Mylar films, 0.7 mm distant from each other. The measurements were carried out with and without the laminar medium to check the contribution of background events. The xenon-scintillator output signals were transmitted through a pulse stretcher and a scintillation-array–controlled linear gate to a 128-channel height analyzer.

In Fig. 2 is given the distribution of the number of events as a function of the energy release in xenon when measured with the laminar medium (closed circles) and without it (open circles). In the first case, the maximum number of events, i.e., the probable value of energy released in

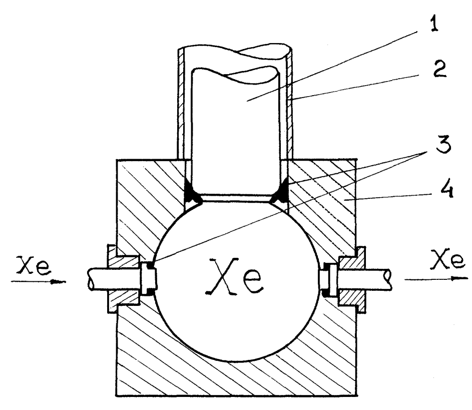

FIG. 1. Schematic of gaseous xenon scintillator. 1, photomultiplier; 2, magnetic shield; 3, Teflon O ring; 4, container.

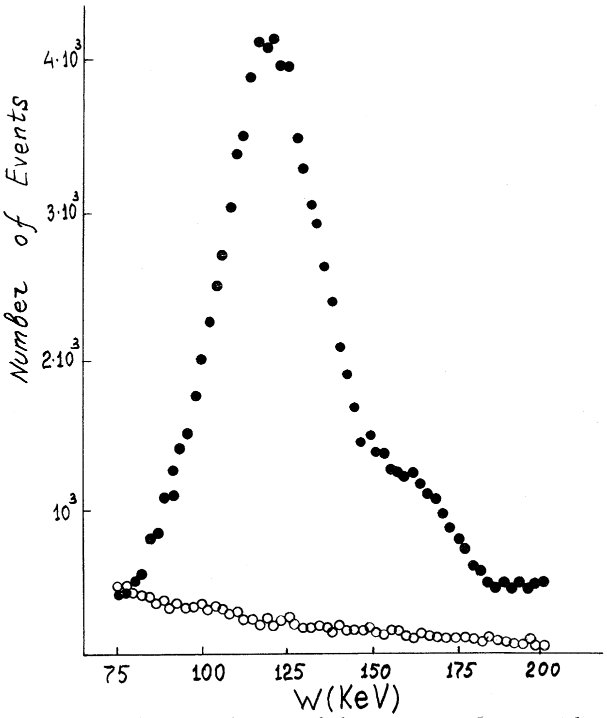

FIG. 2. The distribution of the energy release with the laminar medium (closed circles) and without it (open circles).

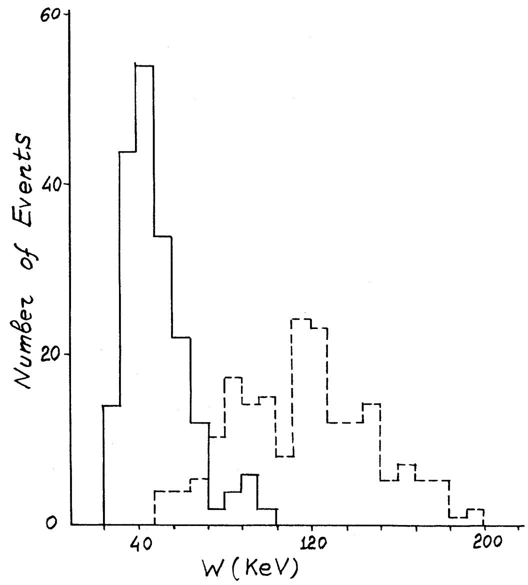

FIG. 3. Monte Carlo distribution of the energy deposited in the xenon scintillator due to ionization losses (dashed) and absorption of XTR and ionization losses (solid).

xenon due to the absorption of XTR quanta and to the ionization losses, corresponds to 125 keV. A calculation of the probable value of ionization losses in xenon gives the figure of 48 keV. Hence, 77 keV is ascribed to the absorption of XTR.

The relation of the total number of events in an interval of energy release 75–200 keV to the number of electrons traversing the laminar medium, i.e., the efficiency of electron registration by their transition radiation, was 0.865 ± 0.095.

The particle registration probability as measured in the absence of the laminar medium was 0.110 ± 0.013. A part of these events is due to the tail of Landau distribution and another part to electron bremsstrahlung in a 12.5-g/cm^2 liquid hydrogen target used in another experiment and installed in front of the laminar medium.

The distribution of 200 events of energy release in xenon, due to XTR and the ionization losses of 31-GeV electrons (Fig. 3, solid curve) as well as to the ionization losses only (Fig. 3, dashed curve), was calculated for our detection instruments by a Monte Carlo technique based on the work of Garibian.[11] A comparison of the curves of Fig. 2 with those of Fig. 3 shows good agreement between experimental results and theoretical predictions.

These data witness that the XTR registration technique by the xenon scintillator allows one to identify untrahigh-energy particles in the region of $\gamma \geqslant 10^3$–10^4.

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Black Holes and Spinning Test Bodies*

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The orbit equations of Papapetrou are the basis for a study of spinning test bodies in a Schwarzschild or Kerr metric field. The constant of the motion for a Killing vector of an arbitrary metric field is given. An analysis of test-body equatorial motion suggests that significant departures from the results of geodesic motion such as larger maximum binding and orbits completely stable against capture may be expected for objects with large intrinsic spin.

Currently much attention is focused on black holes as possible underlying sources for large energy production and/or gravitational wave emission.[1-6] In previous studies of the binding energy available in capture of a test body by a black hole, only geodesic motion has been considered. It is the purpose of this Letter to present results of a study of spinning test bodies which, as is well known, do not execute geodesic motion.

Since a significant fraction of the stars observed have intrinsic angular momentum, a study of a spinning test body in the field of a black hole may serve as a model for the possible star–black-hole interactions. The simple act of endowing a black hole with angular momentum has led to an unexpected richness of possible physical phenomena. It seems appropriate to ask whether endowing the test body with intrinsic spin might not also lead to surprises.[7] Indeed, since in a weak-field approximation the spin-spin interaction is $\propto 1/r^4$ in contrast to the Newtonian attraction which is $\propto 1/r^2$, the possibility exists for spin interactions to dominate completely for small separation.[8,9]

The equations of motion for a spinning test body most commonly used are those of Papapetrou.[10,11] For extended bodies, alternative equations have been suggested with differing supplementary conditions.[12-14] We will, however, confine our attention to a model of point test bodies and use the Papapetrou equations for a spinning point mass as developed by Taub.[15] Our primary interest is in those equations determining the orbit and not the spin equations of motion. There is general agreement, even for extended bodies, that the spin undergoes Fermi-Walker transport and it is just this spin behavior which has been subjected to the most detailed examination.[9,16] Study of the orbit equations has been limited to the Schwarzschild field.[17,18]

Ignoring radiation reaction, the equations describing the motion of a test body with mass m, four-velocity u, and spin four-vector S in a metric field $g_{\mu\nu}$ with curvature operator R are

$$(mu+u\wedge\dot{u}\wedge S)^{\cdot}+\tfrac{1}{2}g^{\rho\mu}R(u,e_\mu)(e_\rho\wedge S\wedge u)=0, \quad (1a)$$

$$\dot{S}\wedge u=0, \quad (1b)$$

$$S\cdot u=0; \quad (1c)$$

e_μ ($\mu=0, \ldots, 3$) denote the basis four-vectors.[19] m and $S\cdot S$ are known to be constants of the motion. It can also be shown that for an arbitrary Killing vector field ξ the scalar

$$(mu+S\wedge u\wedge\dot{u})\cdot\xi+\tfrac{1}{2}g^{\mu\rho}(e_\rho\wedge S\wedge u)\cdot\nabla_\mu\xi \quad (2)$$

is a constant of the motion.[20] The proof is straightforward and relies only on Eq. (1a) and well-known properties of Killing vector fields.

We choose that "specific internal energy" (ϵ of Taub's Eq. [1.5]) to be zero so that the test body is characterized completely by giving its mass and spin. This implies as a consequence that

$$\dot{u}\wedge S=0. \quad (3)$$

Thus, the second term in Eq. (1a) vanishes and one is left with

$$mu\cdot\xi+\tfrac{1}{2}g^{\mu\rho}(e_\rho\wedge S\wedge u)\cdot\nabla_\mu\xi \quad (4)$$