Radiation Bursts from Particles in the Field of Compact, Impenetrable, Astrophysical Objects

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Abstract

The radiation emitted by charged, scalar particles in a Schwarzschild field with maximal acceleration corrections is calculated classically and in the tree approximation of quantum field theory. In both instances the particles emit radiation that has characteristics similar to those of gamma-ray bursters.

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Compact, impenetrable, astrophysical objects (CIAOs)\cite{1, 2, 3, 4} arise in a model, proposed by Caianiello and collaborators\cite{5}, in which particle accelerations have an upper limit $A_m = 2mc^3/\hbar$, referred to as maximal acceleration (MA). The limit can be derived from quantum mechanical considerations \cite{6, 7}, is a basic physical property of all massive particles and must therefore be included in the physical laws from the outset. This requires a modification of the metric structure of space-time.

Classical and quantum arguments supporting the existence of a MA have been discussed in the literature \cite{8, 9, 10}. MA also appears in the context of Weyl space \cite{11}, and of a geometrical analogue of Vigier’s stochastic theory \cite{12} and plays a role in several issues. It is invoked as a tool to rid black hole entropy of ultraviolet divergences \cite{13} and of inconsistencies arising from the application of the point-like concept to relativistic particles \cite{14}. MA may be also regarded as a regularization procedure \cite{15} that avoids the introduction of a fundamental length \cite{16}, thus preserving the continuity of space-time.

An upper limit on the acceleration also exists in string theory where Jeans-like instabilities occur \cite{17, 18} when the acceleration induced by the background gravitational field reaches the critical value $a_c = \lambda^{-1} = (m\alpha)^{-1}$ where $\lambda$, $m$ and $\alpha^{-1}$ are string size, mass and tension. At accelerations larger than $a_c$ the string extremities become casually disconnected. Frolov and Sanchez \cite{19} have also found that a universal critical acceleration must be a general property of strings. It is the same cut-off required by Sanchez in order to regularize the entropy and the free energy of quantum strings \cite{20}.

Applications of Caianiello’s model include cosmology \cite{21}, the dynamics of accelerated strings \cite{22}, the energy spectrum of a uniformly accelerated particle \cite{23}, neutrino oscillations \cite{24, 25}, and photons in a cavity resonator \cite{26}. The model also makes the metric observer–dependent, as conjectured by Gibbons and Hawking \cite{27}.

The consequences of the model for the classical electrodynamics of a particle \cite{28}, the mass of the Higgs boson \cite{29}, and the Lamb shift in hydrogenic atoms \cite{30} have been worked out. In the last instance the agreement between experimental data and MA corrections is very good for $H$ and $D$. For $He^+$ the agreement between theory and experiment is improved by 50% when MA corrections are considered.

MA effects in muonic atoms appear to be measurable \cite{31}. MA also affects the helicity and chirality of particles \cite{32}.

More recently, the model has been applied to particles falling in the gravitational field of a spherically symmetric collapsing object \cite{1}. In this problem MA manifests itself through a spherical shell external to the Schwarzschild horizon and impenetrable to classical and quantum particles \cite{2}. The shell is not a sheer product of the coordinate system, but survives, for instance, in isotropic coordinates. It is also present in the Reissner-Nordström \cite{3} and Kerr \cite{4} cases. The model therefore seems to preclude the usual process of formation of black holes. These are replaced by CIAOs. The purpose of this work is to study the emission of radiation by a scalar particles in the field of a CIAO.

Caianiello’s model is based on an embedding procedure \cite{1} that stipulates that the line element experienced by an accelerating particle is represented by

\begin{equation}
\begin{aligned}
d\tau^2 &= \left(1 + \frac{g_{\mu\nu}\dot{x}^\mu\dot{x}^\nu}{A_{m}^2}\right)g_{\alpha\beta}dx^\alpha dx^\beta \equiv \sigma^2(x)g_{\alpha\beta}dx^\alpha dx^\beta \equiv \bar{g}_{\alpha\beta}dx^\alpha dx^\beta, \\
\end{aligned}
\end{equation}

where $g_{\alpha\beta}$ is a background gravitational field. The effective space-time geometry given by (1)
therefore exhibits mass-dependent corrections that in general induce curvature and violations of the equivalence principle. The four-acceleration $\ddot{x}^\mu = d^2x^\mu/ds^2$ appearing in (1) is a rigorously covariant quantity only for linear coordinate transformations. Its transformation properties are however known and allow the exchange of information among observers. Lack of covariance for $\ddot{x}^\mu$ in $\sigma^2 (x)$ is not therefore fatal in the model. The justification for this choice lies primarily with the quantum mechanical derivation of MA which applies to $\ddot{x}^\mu$, requires the notion of force, is therefore Newtonian in spirit and is fully compatible with special relativity. The choice of $\ddot{x}^\mu$ in (1) is, of course, supported by the weak field approximation (see Figs. 1 and 2 in Ref. [1]). An expansion of (3) in the neighborhood of $\rho = 2$ as $\rho \rightarrow 2$ and $\rho \rightarrow 0$, and that $\ddot{V}_{eff} \rightarrow 1$ as $\rho \rightarrow \infty$. Plots of (3) for different values of $\ddot{E}$ show a characteristic step-like behaviour in the neighborhood of $\rho = 2$ (see Figs. 1 and 2 in Ref. [1]). An expansion of (3) in the neighborhood of $\rho = 2$ yields, in fact, $\ddot{V}_{eff} \sim \ddot{E}^2 + \frac{(\rho-2)^4}{4e^2E^4} + O((\rho-2)^5)$ which has the minimum $\ddot{E}^2$ on the horizon $\rho = 2$. This term vanishes only in the limit $\ddot{E} \rightarrow \infty$ and/or in the limit $\epsilon \rightarrow \infty$, for which $A_m$ or $M$ or both vanish and the problem becomes meaningless.

The addition of MA effects does therefore produce a spherical shell of radius $2 < \rho < 2 + \eta$, with $\eta \ll 1$. The shell is classically impenetrable and remains so at higher orders of approximation [1]. The analogous occurrence of a classically impenetrable shell was first derived by Gasperini as a consequence of the breaking of the local $SO(3,1)$ symmetry [36]. A classically impenetrable shell and shift in horizon also occur in the problem of particles in hyperbolic motion in a Kruskal plane [24].

The existence of large accelerations in proximity of the shell suggests the possibility of radiation of photons (and gravitons) by bremsstrahlung. It has become recently clear that
production of electromagnetic radiation is possible even when the acceleration is produced by a gravitational field \[ \Box, \Box, \Box, \Box, \Box, \Box. \] Two approaches are followed here.

i) Generalization of Larmor’s formula. A covariant generalization of this formula would replace \( \dot{x}^\mu \) with its covariant counterpart which vanishes rigorously along a geodesic in general relativity. It is of course known that even in Einstein’s theory a charged particle, while trying to adhere to the equivalence principle, does not really follow a geodesic in a gravitational field. Deviations necessarily develop an account of damping and of the tail effect discussed by De Witt and Brehme \[ \Box. \] In Caianiello’s model, not only is the equivalence principle violated, but the MA corrections also introduce deviations from geodetic motion. These are introduced by \( \sigma^2(x) \) which in first approximation is given by \( \Box \). From the definition \( \Box \) of \( \sigma^2(x) \) one finds \( g_{\mu \nu} \dot{x}^\mu \dot{x}^\nu = (\sigma^2 - 1)A^2_{\text{eff}} \) and Larmor’s formula, which applies to particles of arbitrary trajectory, therefore becomes

\[
P \simeq -\frac{2q^2}{3} A^2_{\text{eff}} (\sigma^2 - 1). \tag{4}
\]

When the MA corrections vanish, \( \sigma \to 1 \) and \( P \to 0 \). Eq. \( \Box \) therefore seems suitable to deal with radiation due to MA. On the basis of \( \Box \) and \( \Box \) one may expect radiation to be particularly intense when \( \rho \) is close to 2, where \( \sigma^2(x) \) has a singularity. One finds in fact

\[
P \simeq -\frac{2q^2}{3} A^2_{\text{eff}} \frac{\epsilon^2}{\rho^7 (\rho - 2)^3} \left\{ \lambda^4 (\rho - 2)^2 (7 - 10 \rho + 3 \rho^2) + \rho^4 [-4 + 4(1 + 2 E^2) \rho + \rho^2 (1 - 4 E^2 + 4 E^4)] - 2 \lambda^2 \rho^2 (\rho - 2) [10 + (-19 + 10 E^2) \rho + (11 - 8 E^2) \rho^2 + 2(\bar{E}^2 - 1) \rho^3] \right\}. \tag{5}
\]

However, the particle can not reach the value \( \rho = 2 \) because of its acceleration and ensuing impenetrable shell. The shell’s radius corresponds to a maximum of \( \bar{V}_{\text{eff}}^2 \) and to the smallest value of \( \rho \) the particle can reach. By expanding \( \bar{V}_{\text{eff}}^2 \) as a series in the neighborhood of \( \rho = 2 \) and keeping terms to sixth order, one finds

\[
\bar{V}_{\text{eff}}^2 \sim \left\{ 256 \bar{E}^2 (\rho - 2)^4 \epsilon^2 + 5(4 + \lambda^2)^2 (\rho - 2)^6 \epsilon^2 + 1024 \bar{E}^{10} \epsilon^4 - 16 \bar{E}^2 (\rho - 2)^5 \times \cdot \cdot \cdot \right\} \frac{1}{1024 \bar{E}^8 \epsilon^4}, \tag{6}
\]

whose maximum is at

\[
\rho_1 = \left\{ 2(768 \bar{E}^2 - 240 \epsilon^2 + 272 \bar{E}^2 \epsilon^2 - 120 \lambda^2 \epsilon^2 + 164 \bar{E}^2 \lambda^2 \epsilon^2 - 15 \lambda^4 \epsilon^2) \right. \tag{7}\]
\[
+ 4 \bar{E}^2 \epsilon \sqrt{6144 \bar{E}^2 - 1520 \epsilon^2 + 1536 \bar{E}^2 \epsilon^2 - 760 \lambda^2 \epsilon^2 + 1152 \bar{E}^2 \lambda^2 \epsilon^2 - 95 \lambda^4 \epsilon^2} \right\} \times \left\{ 3(256 \bar{E}^2 - 80 \epsilon^2 + 64 \bar{E}^2 \epsilon^2 - 40 \lambda^2 \epsilon^2 + 48 \bar{E}^2 \lambda^2 \epsilon^2 - 5 \lambda^4 \epsilon^2) \right\}^{-1}.
\]

The distance of closest approach \( \rho_1 \) is real for

\[
\bar{E}^2 \geq \epsilon^2 (1520 + 760 \lambda^2 + 95 \lambda^4) (6144 + 153 \epsilon^2 + 1152 \lambda^2 \epsilon^2)^{-1}
\]

and always positive for small values of \( \lambda \) (radial infall) and \( \epsilon \) which are appropriate in many instances. It is useful to recall at this point the relationship between the adimensional parameters \( M, \bar{E}, \lambda, \epsilon, \tau \) (adimensional time) and the corresponding (primed) CGS quantities. One has

\[
M = \frac{GM'}{c^2}, \quad \bar{E} = \frac{E'}{m'c^2}, \quad \lambda = \frac{L'c}{GM'm'}, \quad \epsilon = \frac{hc}{2GM'm'}, \quad \tau = \frac{c^3t'}{GM'}.
\]
If $M' \sim 3M_\odot$ and the falling particle is a proton of energy $E' \sim 10\text{MeV}$, one finds $\epsilon \sim 4.7 \times 10^{-20}$ and $\tilde{E} \sim 0.01$. For these values of the parameters $2 - \rho_1 \simeq 9.7 \times 10^{-16}$.

An expression that gives $P$ as a function of $M$, $\tilde{E}$, $\lambda$ and $\epsilon$ can be obtained by substituting (7) into (5). Its usefulness is however limited because the parameters involved can not yet be linked directly to the particle’s motion in the new field $\tilde{g}_{\mu\nu}$. This can be achieved by writing

$$\frac{dr}{d\tilde{s}} = \frac{dx^0}{d\tilde{s}} = \frac{\tilde{E}}{\sigma^2 g_{00}}$$

and using

$$\left(\frac{dr}{d\tilde{s}}\right)^2 = \frac{\tilde{E}^2 - \tilde{V}_{\text{eff}}^2}{\tilde{V}_{\text{eff}}^2}$$

The desired function $r = r(x^0)$ is obtained by integrating (8) with respect to time. In dimensional variables, one finds

$$\tilde{E} \int \frac{\rho d\rho}{\sigma^2 (\rho - 2) \sqrt{\tilde{E}^2 - \tilde{V}_{\text{eff}}^2(\rho)}} = -\tau + a$$

where the negative sign in front of $\tau$ accounts for the fact that the particle approaches the gravitational source as $\tau$ increases. The constant $a$ must be determined by appropriate initial conditions.

By assuming that the motion is radial ($\lambda = 0$) and for the values of $\tilde{E}$ and $\epsilon$ given above, the l.h.s. of (9) can be integrated in the neighborhood of $\rho_1$ and yields

$$2\sqrt{2} \ln(\rho - \rho_1) \sim -\tau + a.$$  

By requiring that $\rho \simeq 2.1$ at $t = 0$, one finds $a = -1.3272$ and (10) becomes

$$\rho \simeq \rho_1 + e^{-(\tau + 0.5127)/2\sqrt{2}}.$$  

The particle therefore reaches a distance within a 1% of $\rho_1$, in the time $t' \geq 4.8 \times 10^{-5}\text{s}$. However, radiation becomes particularly intense when $t' \geq 1.68 \times 10^{-3}$. If (11) is then substituted into (8), one obtains

$$|P| \sim \frac{2q^2 A_m^2}{3GM'}\left[2.67 \times 10^{42} + 3.85 \times 10^{47}(t' - 1.68 \times 10^{-3}) + 3.78 \times 10^{52}(t' - 1.68 \times 10^{-3})^2\right].$$

Radiation becomes appreciable $\sim 10^{30}\text{erg/s}$ only at $t' \geq 5 \times 10^{-4}\text{s}$ and increases very rapidly with $t'$. As shown in Fig. 1, it already is $P \sim 10^{51}\text{erg/s}$ at $t' \sim 10^{-3}\text{s}$. Because of the quadratic dependence in both particle charge and mass, small clumps of matter radiate significantly more and at higher frequencies even a certain distance from $\rho_1$.

**ii) Tree-level calculations.** The second way of estimating the power produced follows the procedure of Crispino, Higuchi and Matsas [41]. It treats the electromagnetic field as a scalar field, whose source is the current produced by a particle, also a scalar. Quantum field theory at the tree level is then used to describe the electromagnetic field in a Schwarzschild background. The method is particularly interesting in the context of Caianiello’s model because it enables a comparison between the massive scalar particle that produces the current and experiences the
effective geometry $\sigma^2g_{\mu\nu}$, and the massless photon that can not be accelerated and therefore "sees" only a Schwarzschild geometry. The source is the scalar current

$$j(x) = q \sqrt{-g^0} \delta(r - r_S)\delta(\theta - \pi/2)\delta(\phi),$$

(13)

normalized by the requirement that $\int d\sigma_V j(x) = q$, where $d\sigma_V$ is the proper 3-volume element orthogonal to $\bar{u}^\mu$, the four-velocity of the source. The four-velocity component $\bar{u}^0 = \bar{E}g^{00} = \frac{\bar{E}}{\sigma\sigma_0}$ takes into account the MA corrections to the motion of the scalar particle, whereas the 3-volume integration is in Schwarzschild space, hence the $\sqrt{-g}$ in (13). The scalar electromagnetic field generated by (13) is, at the tree level,

$$A_{\omega lp} = i \int d^4x \sqrt{-g} j(x) u_{\omega lp}^*,$$

(14)

where the functions

$$u_{\omega lp}(x) = \sqrt{\frac{\omega}{\pi}} R_{\omega l}(r) Y_{l\ell}(\theta, \phi)e^{-i\omega t}$$

(15)

represent the positive-frequency solutions ($\omega > 0$) in Schwarzschild space-time of the equation $\nabla_\mu \nabla_\mu u_{\omega lp} = 0$. The functions $R_{\omega l}$ satisfy the equation

$$e^{2\lambda} R''_{\omega l} + \left(\frac{2}{r} + \lambda'\right) e^{2\lambda} R'_{\omega l} + \left[\omega^2 - e^{\lambda}\frac{l(l + 1)}{r^2}\right] R_{\omega l} = 0,$$

(16)

where $e^\lambda = 1 - 2M/r$ and the primes denote differentiation with respect to $r$ [2]. The emitted power is then [41]

$$P = \int_0^\infty d\omega \frac{|A_{\omega lp}|^2}{T},$$

(17)

where $T$ is the time during which the interaction is switched on [14]. Introducing (13) into (14) and using the coordinate $\rho - 2 = x$, one finds, to leading order,

$$A_{\omega lp} \approx \frac{D e^{6.512/\sqrt{\omega}} e^{iM\omega + 1/\sqrt{\omega}}}{2M \sqrt{\omega}} Y_{l\ell}(\pi/2, 0)(1 + \bar{R}_{\omega l}),$$

(18)

where

$$D \equiv i \frac{4q\bar{E}^3}{A_{\omega lp}^2 M \sqrt{\pi}} Y_{l\ell}(\pi/2, 0)(1 + \bar{R}_{\omega l}).$$

The contribution of the term $-1$ on the r.h.s. of (18) is negligible relative to the exponential. The reflection coefficient $\bar{R}_{\omega l} < 1$ is also neglected.

Eqs. (17) and (18) give the radiated power as

$$|P| \approx \frac{q^2 \bar{E}^6 e^{6.512} e^{\sqrt{\omega}}}{\pi A_{\omega lp}^2 M^6} \frac{e^{\sqrt{\omega} \tau}}{\tau},$$

(19)

where the time $T = M\tau'$ in (17) and $\tau'$ is the time required for the particle to reach the intense radiation zone of its trajectory. Though this result vanishes as $\hbar \to 0$, the correct limit for
vanishing MA contributions follows by taking the limit $\sigma \to 1$ in (13) and is finite. In CGS units, the power becomes

$$|P| \simeq \frac{q^2 E^6 h^4 e^{8.21} e^{\sqrt{2} c^3 t'/GM'}}{16\pi m' G^5 M'^5} \frac{e^{\sqrt{2} c^3 t'/GM'}}{t'} \simeq 2.35 \times 10^{-110} e^{1.9 \times 10^5 t'/t'}. \quad (20)$$

Despite the enormously small reducing factor in (20), a single proton reaches powers of $\sim 10^{30}\text{erg/s}$ after only $t' \sim 2 \times 10^{-3}\text{s}$, and $P \sim 10^{51}\text{erg/s}$ is reached for $t' \sim 2.3 \times 10^{-3}\text{s}$, in agreement with the results of (12). The frequency distribution of $P$ is given by

$$\frac{dP}{d\omega} \simeq \frac{|D|^2 e^{9.121}}{4M^3} \frac{e^{\sqrt{2} \tau}}{(M^2 \omega^2 + \frac{1}{2})^2}. \quad (21)$$

or, in CGS units, by

$$\frac{dP}{d\omega} \simeq \frac{10^4 h^4 E^6 q^2}{4\pi^2 G^4 M'^4 \alpha^{10} c^8} \frac{e^{\sqrt{2} c^3 t'/GM'}}{(G^2 M'^2 \omega^2 + \frac{1}{2})^2 t'}. \quad (22)$$

It can be seen from the Figs. 2 and 3 that the bursts of radiation are sudden and peaked at lower frequencies, though the intensity of radiation also remains high at $\gamma$-ray frequencies. Some similarities with $\gamma$-ray busters are obvious. The latter have extraordinary large energy outputs $\sim 10^{51}$ to $\sim 10^{54}\text{erg/s}$. Their spectra are non-thermal with photon energies ranging from $50\text{keV}$ to the $\text{GeV}$ region. In CIAO’s case spectra also are non-thermal, in fact typically of the bremsstrahlung type, with intense radiation that extends to the same frequency range of bursters. Significant amounts of low frequency radiation could be absorbed by the progenitor, if this is a black hole, in the case of bursters. Black holes are known to have sizeable absorption for infrared photons [41]. For CIAOs absorption could also take place at the shell which, for photons, behaves like a horizon. Nonetheless energy extraction from CIAOs remains extremely efficient. The duration of the intense emission by particles in their proximity is of the same order of magnitude, $1\text{ms}$ to $10^3\text{s}$, of the events that are observed. Finally, the lower frequency afterglows that are sometimes observed to last several months, may be explained, in CIAO’s scenario, as due to a dip in the effective potential experienced by scalar particles at higher accelerations[2]. The material trapped in this way may subsequently leave the relatively shallow attractive area to resume the acceleration and radiation processes at lower intensities and frequencies.
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References

[1] A. Feoli, G. Lambiase, G. Papini, G. Scarpetta, Phys. Lett. A 263 (1999) 147.
[2] S. Capozziello, A. Feoli, G. Lambiase, G. Papini, G. Scarpetta, Phys. Lett. A 268 (2000) 247.
[3] V. Bozza, A. Feoli, G. Papini, G. Scarpetta, Phys. Lett. A 271 (2000) 35.
[4] V. Bozza, A. Feoli, G. Lambiase, G. Papini, G. Scarpetta, Phys. Lett. A 283 (2001) 53.
[5] E.R. Caianiello, La Rivista del Nuovo Cimento 15 n.4 (1992) and references therein.
[6] E.R. Caianiello, Lett. Nuovo Cimento 41 (1984) 370.
[7] W.R. Wood, G. Papini and Y.Q. Cai, Il Nuovo Cimento 104B, 361 and (errata corrige) 727 (1989).
[8] A. Das, J. Math. Phys. 21 (1980) 1506; M. Gasperini, Astrophys. Space Sci. 138 (1987) 387; M. Toller, Nuovo Cimento 102B (1988) 261; Int. J. Theor. Phys. 29 (1990) 963; Phys. Lett. B 256 (1991) 215; B. Mashhoon, Physics Letters A 143 (1990) 176; V. de Sabbata, C. Sivaram, Astrophys. Space Sci. 176 (1991) 145; ”Spin and Torsion in gravitation”, World Scientific, Singapore, (1994); D.F. Falla, P.T. Landsberg, Il Nuovo Cimento 106B, (1991) 669; A.K. Pati, Il Nuovo Cimento 107B (1992) 895; Europhys. Lett. 18 (1992) 285.
[9] C.W. Misner, K.S. Thorne, J. A. Wheeler, "Gravitation", W.H. Freeman and Company, S. Francisco, 1973.
[10] H.E. Brandt, Lett. Nuovo Cimento 38, (1983) 522; Found. Phys. Lett. 2 (1989) 3.
[11] G. Papini and W.R. Wood, Phys. Lett. A 170 (1992) 409; W.R. Wood and G. Papini, Phys. Rev. D 45 (1992) 3617; Found. Phys. Lett. 6 (1993) 409; G. Papini, Mathematica Japonica 41 (1995) 81.
[12] J. P. Vigier, Found. Phys. 21 (1991) 125.
[13] M. McGuigan, Phys. Rev. D 50 (1994) 5225.
[14] G.C. Hegerfeldt, Phys. Rev. D 10 (1974) 3320.
[15] V.V. Nesterenko, A. Feoli, G. Lambiase, G. Scarpetta, Phys. Rev. D 60 (1999) 965001.
[16] See, for instance: J.C. Breckenridge, V. Elias, T.G. Steele, Class. Quantum Grav. 12 (1995) 637.

[17] N. Sanchez and G. Veneziano, Nucl. Phys. B 333 (1990) 253; M. Gasperini, N. Sanchez, G. Veneziano, Nucl. Phys. B, 364 (1991) 365; Int. J. Mod. Phys. A 6 (1991) 3853.

[18] M. Gasperini, Phys. Lett. B 258 (1991) 70; Gen. Rel. Grav. 24 (1992) 219.

[19] V.P. Frolov and N. Sanchez. Nucl. Phys. B 349 (1991) 815.

[20] N. Sanchez, in “Structure: from Physics to General Systems” eds. M. Marinaro and G. Scarpetta (World Scientific, Singapore, 1993) vol. 1, pag. 118.

[21] E.R. Caianiello, M. Gasperini, G. Scarpetta, Class. Quantum Grav. 8 (1991) 659; M. Gasperini, in “Advances in Theoretical Physics” ed. E.R. Caianiello, (World Scientific, Singapore, 1991), p. 77.

[22] A. Feoli, Nucl. Phys. B 396 (1993) 261.

[23] E.R. Caianiello, A. Feoli, M. Gasperini, G. Scarpetta, Int. J. Theor. Phys. 29 (1990) 131.

[24] E.R. Caianiello, M. Gasperini, and G. Scarpetta, Il Nuovo Cimento 105B (1990) 259.

[25] V. Bozza, G. Lambiase, G. Papini, G. Scarpetta, Phys. Lett. A 279 (2001) 163. V. Bozza, S. Capozziello, G. Lambiase, G. Scarpetta, Int. J. Theor. Phys. 40 (2001) 849.

[26] G. Papini, A. Feoli, and G. Scarpetta, Phys. Lett. A 202 (1995) 50.

[27] G. Gibbons and S.W. Hawking, Phys. Rev. D15 (1977) 2738.

[28] A. Feoli, G. Lambiase, G. Papini, and G. Scarpetta, Il Nuovo Cimento 112B (1997) 913.

[29] G. Lambiase, G. Papini, and G. Scarpetta, Il Nuovo Cimento 114B (1999) 189. See also S. Kuwata, Il Nuovo Cimento 111B (1996) 893.

[30] G. Lambiase, G. Papini, and G. Scarpetta, Phys. Lett. A 244 (1998) 349.

[31] C.X. Chen, G. Lambiase, G. Mobed, G. Papini, and G. Scarpetta, Il Nuovo Cimento 114B (1999) 1135.

[32] C.X. Chen, G. Lambiase, G. Mobed, G. Papini, and G. Scarpetta, Il Nuovo Cimento 114B (1999) 199.

[33] C. Lämmerzahl, Gen. Rel. Grav. 28(1996)1043.

[34] D. Singh and G. Papini, Il Nuovo Cimento 115B(2000)223.

[35] S. Weinberg, ”Gravitation and Cosmology”, John Wiley and Sons, New York, 1972, Ch.4.

[36] M. Gasperini, Phys. Rev. D34 (1986) 2260.

[37] David G. Boulware, Ann. Phys. (N.Y.) 124 (1980) 169.
[38] A. Higuchi, G.E.A. Matsas and D. Sudarsky, Phys. Rev. D 46 (1992) 3450.
[39] Jürgen Audretsch, Ulf Jasper, Vladimir D. Sharzhinsky, Phys. Rev. D 49 (1994) 6576.
[40] Atsushi Higuchi, George E. A. Matsas, Daniel Sudarsky, Phys. Rev. D 56 (1997) R6071.
[41] L.C.B. Crispino, A. Higuchi, G.E.A. Matsas, Class. Quant. Grav. 17 (2000) 19.
[42] Luis C. B. Crispino, Atsushi Higuchi, George E.A. Matsas, Phys. Rev. D 63 (2001) 124008.
[43] B. We Witt, R.W. Brehme, Ann. Phys. (N.Y.) 9 (1960) 221.
[44] C. Itzykson and J.-B. Zuber, ”Quantum Field Theory”, McGraw-Hill, New York, 1980.
Figure 1: Power emitted as a function of time for $M' = 3M_\odot$, $E' = 10\text{MeV}$, $m' = 1.7 \times 10^{-24} \text{g}$

Figure 2: Power emitted as a function of time according to Eq. (20). The values of $M'$, $E'$, $m'$ are as in Fig. 1

Figure 3: Power spectrum in the $\gamma$-ray frequency region and the time interval 2.1ms to 2.3ms. The values of $M'$, $E'$, $m'$ are as in Fig. 1