Inverse Compton scattering off black body radiation in high energy physics and gamma (MeV – TeV) astrophysics

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1. Introduction

Inverse Compton Scattering (ICS) plays a relevant role in very high energy astrophysics (cosmic rays and gamma astronomy) [1 – 6] and high energy physics (LEP I, LEP II, accelerators) [7 – 9]. Indeed on the one hand the ICS of high relativistic cosmic rays (either electron at GeV or proton and nuclei at much higher energies) off electromagnetic fields (either cosmological Black Body Radiation (BBR) at \( T \approx 2.73 \) K, interstellar lights, radio waves or even stationary magnetic fields) is a source of high energy photons (X, gamma rays) which we observe in the Universe as diffuse or point source, on the other hand ICS is often the main process responsible for slowing down (i.e. for the energy losses) energetic charged particles; indeed ICS is often the main factor influencing cosmic ray lifetimes as well as their energy spectrum depletion in the harder regions (i.e. of their detailed spectrum shape and evolution) [5, 6]. Moreover, ICS of relativistic jets of compact objects off thermal photons of a stellar companion in a binary system or accretion disk might be, as recently proposed [10], the key process able to produce ‘gamma jets’ responsible (by their rotation and blazing into different directions) for the puzzling Gamma Ray Bursts (GRB). So even if this subject seems settled [3, 4] we revisit the ICS process in order to obtain an analytic and compact formula able to describe the differential energy and angular spectra of relativistic charges resulting from ICS off a BBR. One of the main features of our final expressions is that we can easily cover the entire range of energy of the ICS energy spectrum. These results improve the Monte Carlo simulation because the latter must consider only a thin portion of the energy range in order to inspect the ICS spectrum in detail [9]. In the following we describe our approach and we show some remarkable results.

2. The spectrum of ICS off BBR

We follow a standard procedure to get the general ICS spectrum of a relativistic charge (electron, proton, nuclei in cosmic rays or in accelerators bunches) hitting photons whose distribution is a BBR spectrum. We consider first the photon target distribution in the Laboratory Frame (LF) where the BBR is isotropic and homogeneous, then we transform it to the Electron Frame (EF) where the BBR is still homogeneous but highly anisotropic, therefore in the EF we evaluate the usual Compton scattering and finally we transform the diffused differential photon number back to the LF. The starting photon target distribution, in the LF, is the well-known isotropic and homogeneous BBR whose number density per unit energy \( dN_0/dE_0 \) and solid angle \( d\Omega_0 \) is given by the Planck formula

\[
\frac{dN_0}{dE_0 d\Omega_0} = \frac{2}{(\pi \hbar)^3} \left\{ \exp\left(\frac{E_0}{k_B T}\right) - 1 \right\}. \tag{1}
\]

where \( k_B \) is the Boltzmann constant and \( T \) the BBR temperature. We transform this distribution to the EF by

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standard Lorentz boosts choosing the $z$ axis as the direction coincident with the initial electron momentum and recalling that $dE/dz$ is a relativistic invariant [2]. In the following we label the quantities related to the electron frame EF by a $(\ast)$, with a subscript 0 if they are considered before the scattering, or subscript 1 if after; so we have
\[
\cos \theta_0 = \frac{\cos \theta_0^* + \beta}{1 + \beta \cos \theta_0^*}, \quad \phi_0 = \phi_0^*, \quad \epsilon_0 = \gamma \epsilon_0^*(1 + \beta \cos \theta_0^*),
\]
where $\beta$ is the dimensionless electron velocity and $\gamma$ the corresponding Lorentz factor. The transformed BBR number density distribution exhibits, in this relativistic limit and in the EF, a clear dipole anisotropy and becomes more and more peaked around $\theta_0^* = \pi$ as $\gamma$ increases. An analogous dipole signature, in the non-relativistic regime, is the cosmological one found at the millikelvin level in the $T = 2.73$ K BBR due to the Earth's motion. The associated energy spectrum is mostly 'green shifted' showing a maximum around $\gamma \epsilon_0 T$ with respect to the original 'red' BBR spectrum with a maximum around $\epsilon_0 T$. The next step is to derive the total number of diffused photons in the EF; this number can be obtained as follows:
\[
\frac{dN_1}{d\epsilon_1^*} = \frac{d\sigma_C}{d\epsilon_1^*} (d\epsilon_1^* d\Omega_1^* d\Omega_0^*) = \frac{d\sigma_C}{d\epsilon_1^*} d\epsilon_1^* d\Omega_1^* c,
\]
where $d\sigma_C/(d\epsilon_1^* d\Omega_1^*)$ is the Compton differential cross section
\[
\frac{d\sigma_C}{d\epsilon_1^* d\Omega_1^*} = \frac{\epsilon_0^*}{2} \left( \frac{\epsilon_1^*}{\epsilon_0^*} \right)^2 \left( \epsilon_1^* + \epsilon_0^* - \sin^2 \theta_0^* \right) \delta \left[ \epsilon_1^* - \left( \frac{\epsilon_0^*}{1 - \epsilon_0^* (1 - \cos \theta_0^*)/mc^2} \right) \right].
\]
Equation (4) is the last step to derive the exact ICS differential number distribution in the Laboratory Frame (where the scattered photons will be observed) by using the inverse Lorentz transformations; the result is
\[
\frac{dN_1}{d\epsilon_1^* d\Omega_1^* d\Omega_0^*} = \frac{\epsilon_1^*}{c\hbar^3} \left[ \begin{array}{c} \gamma \epsilon_1 (1 - \beta \cos \theta_1) (1 + \beta \cos \theta_1^*) \\ \gamma \epsilon_1 (1 - \beta \cos \theta_1) (1 + \beta \cos \theta_1^*) \\ - \gamma \epsilon_1 (1 - \beta \cos \theta_1) (1 - \cos \theta_0^*) + \cos^2 \theta_0^* \\ \end{array} \right] \times \frac{d\Omega_0^*}{\gamma \epsilon_1 (1 - \beta \cos \theta_1) (1 - \cos \theta_0^*)},
\]
where from Eqn (4) $\cos \theta_0^*$ must be expressed as a function of $\theta_0^*, \theta_1^*, \phi_0^*, \phi_1^*$. This ‘blue shifted’ distribution is different from the previous ones because now the expected peak around a given energy $\epsilon_1$ has been spread into a wide plateau from $\epsilon_0 T$ up to $\gamma \epsilon_0 T$. The most general ICS differential distribution can be obtained by means of numerical integration of Eqn (5) over $\Omega_0^*$ but as we are more interested in high energy phenomena we also show how to simplify the above formula in these cases. First of all we recall that in particle accelerators we are dealing with ultrarelativistic particles, i.e. $\gamma \gg 1$; moreover for most astrophysical problems the photon energy in the EF is much smaller than the electron rest mass, i.e. $\epsilon_0^* \ll mc^2$; so it is often possible to approximate the Compton differential cross section by the Thomson one; as a third further approximation we may consider the ultrarelativistic-Thomson limit where the two previous conditions $\gamma \gg 1$ and $\epsilon_0^* \ll mc^2$ are both satisfied. Let us discuss these three different expansions. In the first case ($\gamma \gg 1$) the BBR photons, in the EF, are practically all incident head-on so the incidence angle $\theta_0^*$ can be approximately written as $\theta_0^* \approx \pi - 1/\gamma$. Consequently the scattering angle is related to the $\theta_1^*$ angle by the simple formula $\theta_0^* + \theta_1^* \approx \pi$ and $\cos \theta_0^* \approx -\cos \theta_1^*$ within $1/\gamma$. Moreover the kinematics of the ICS shows that the scattered radiation, in the EF, is strongly concentrated in a narrow cone $\theta_1 \approx 1/\gamma$ along the direction of motion of the ultrarelativistic particle. The ICS differential distribution in Eqn (5), in the ultrarelativistic expansion, reduces to an analytical expression
\[
\frac{dN_1}{d\epsilon_1^* d\Omega_1^* d\Omega_0^*} = \frac{2\pi \epsilon_0 T \gamma \epsilon_1}{c^3 \hbar^3} \times \epsilon_1 \left\{ \begin{array}{l} \left[ 1 - \exp\left( -2\gamma \epsilon_1 (1 - \beta \cos \theta_1)/(2mc^2) \right) \right] \\ \times \frac{\epsilon_1}{\left[ 1 - \epsilon_1 (1 + \cos \theta_1)/(2mc^2) \right]^2} \\ \times \left[ 1 - \epsilon_1 (1 + \cos \theta_1)/(2mc^2) \right]^{-1} \left[ 1 - \epsilon_1 (1 + \cos \theta_1)/(2mc^2) \right]^{-1} \right. \\
\left. \left. - \epsilon_1 (1 + \cos \theta_1)/(2mc^2) + \left( \cos \theta_1 - \beta \cos \theta_1 \right)^2 \right] \right\},
\]
where $1 - \beta \approx 1/(2\gamma^2)$ and $d\Omega_1 \approx 2\pi \theta_0 d\theta_1$. The second possible approximation is the Thomson limit $\epsilon_0^* \ll mc^2$. In this case we can neglect all terms of order $\epsilon_0^*/mc^2$ with respect to 1 in Eqn (5) and the resulting differential distribution becomes
\[
\frac{dN_1}{d\epsilon_1^* d\Omega_1^* d\Omega_0^*} = \frac{2\pi \epsilon_0 T \gamma \epsilon_1}{c^3 \hbar^3} \times \epsilon_1 \left\{ \begin{array}{l} \left[ 1 - \exp\left( -2\gamma \epsilon_1 (1 - \beta \cos \theta_1)/(1 + \beta \cos \theta_1)/(2\gamma \epsilon_0 T) \right) \right] \\ \times \frac{\epsilon_1}{\left[ 1 - \epsilon_1 (1 + \cos \theta_1)/(1 + \beta \cos \theta_1)/(2\gamma \epsilon_0 T) \right]^2} \\ \times \left[ 1 + \left( \cos \theta_1 - \beta \cos \theta_1 \right)^2 \right] + \left[ 1 - \epsilon_1 (1 + \cos \theta_1)/(1 + \beta \cos \theta_1)/(2\gamma \epsilon_0 T) \right]^{-1} \left[ 1 - \epsilon_1 (1 + \cos \theta_1)/(1 + \beta \cos \theta_1)/(2\gamma \epsilon_0 T) \right]^{-1} \right. \\
\left. \left. \times \sin \theta_0^* \cos \theta_0^* d\theta_0^* \right] \right\}.
\]
Indeed, from the previous Eqns (6), (7), we can get the last expansion, the Thomson-ultrarelativistic formula. It stems
from assuming $\gamma \gg 1$ and $\gamma' \ll mc^2$ at the same time. This means that the ICS differential distribution can be obtained by setting $\epsilon_1 \ll mc^2$ in Eqn (6) and by setting $\gamma \gg 1$ in Eqn (7). In the first derivation we obtain:

$$\frac{dN_1}{d\epsilon_1 d\Omega_1} \propto \epsilon_1 \ln \left\{ \frac{1 - \exp \left[ -2\gamma' \epsilon_1 \left( 1 - \beta \cos \theta_1 \right) / (x_0 T) \right]}{1 - \exp \left[ -\epsilon_1 \left( 1 - \beta \cos \theta_1 \right) / (2x_0 T) \right]} \right\} \times \epsilon_1 \ln \left[ 1 + \left( \frac{\cos \theta_1 - \beta}{1 - \beta \cos \theta_1} \right)^2 \right]. \quad (8)$$

In the second derivation the integral contained in Eqn (7) can be simplified remembering that $\theta_0 \approx \pi - (1/\gamma)$ so $\sin \theta_0 \cos \theta_0 \approx 1/\gamma$ and the integrand is significantly different from zero only inside a thin cone whose aperture is just $1/\gamma$. Thus the result is a formula differing from Eqn (8) only by a factor of order $1/\gamma^2$ and in our assumptions this term is completely negligible. So the still analytical Eqn (8) is our ultrarelativistic-Thomas distribution. In Figure 1 we show the surprising metamorphosis of a Planckian Black Body into a taller (by a suppression factor $\gamma^2$) serial of smooth truncated hills at larger relativistic regimes. This process is mainly due to the overlapping of a series of blueshifted spectra obtained at different angular directions; the softer photons are the more isotropic ones while the harder photons are strongly anisotropic and located in the inner cone of the beam. The ‘Rayleigh’ regions of the differential ‘BBR’ spectra overlap each other even if calculated at different $\theta_1$ angles ($0 < \theta_1 < \pi$) while their peaks are higher and more and more blueshifted for $\theta_1$ angles approaching zero. Consequently the exponential decay of the ICS spectrum on the right side reflects the ‘Wien’ behavior regions of the angular ‘BBR’ spectra at $\theta_1 \approx 0$. We give here a summarized characteristic of the ICS energy spectrum behavior over the whole $\epsilon_1$ energy range. For $\epsilon_1 \ll x_0 T / \gamma'$ the spectrum exhibits linear growth (like the Rayleigh region of a BBR one); for $x_0 T / \gamma' < \epsilon_1 < x_0 T$ the linear behavior is modified by a logarithmic correction and the growth is proportional to $\epsilon_1 \ln (4\gamma' x_0 T / \epsilon_1)$; for $x_0 T < \epsilon_1 < 4\gamma' x_0 T$ the spectrum is spread into a very slowly linearly decreasing plateau up to $\epsilon_1 \approx 4\gamma' x_0 T$; for $\epsilon > 4\gamma' x_0 T$ the spectrum decays as $\epsilon_1 \exp(-\epsilon_1 / 4\gamma' x_0 T)$ [we remind the reader that the Wien region of a BBR distribution decays as $\epsilon_1^2 \exp(-\epsilon_1 / x_0 T)$].

We recall that the commonly and widely applied analytical formulae on ICS are those derived by F Jones [3] and based on ICS off monochromatic and isotropic radiation; this factitious and artificial ‘BBR’ leads to a final spectrum in disagreement with those experimentally observed. In the next two sections we show that we can correctly evaluate ICS spectra in high energy physics and astrophysics by means of our formulae.

3. The ICS spectrum in LEP experiments

We could directly verify the validity of present results by a (successful) fit of experimental ICS spectra obtained at LEP by the A Melissinos group [7] and by the G Diambrini-Palazzi group [8] (at a higher degree of precision). Indeed the LEP vacuum pipe can be considered as a black body cavity at room temperature ($T \approx 291$ K); hence the electromagnetic radiation in thermal equilibrium is scattered and beamed ahead by electron (and positron) bunches whose energy is $E_0 = 45.6$ GeV. This energy corresponds to a Lorentz factor $\gamma = 8.92 \times 10^4$ so we can apply our ultrarelativistic formulae for ICS to evaluate the interesting region of the spectrum. We compare our results with the Monte Carlo simulations performed at LEP and fit the ICS spectrum to the experimental data. We compute the Thomson and the Compton spectra in order to show the differences between the two curves and with respect to the Monte Carlo simulation. The spectra can be obtained by numerical integration over the $\theta_1$ angle in expressions (6)–(8) respectively for the Compton or Thomas limit:

$$\frac{dN_{1(T,K)}}{d\epsilon_1} = N_0 \int_{D_0} \frac{dN_{1(T,K)}}{d\epsilon_1 d\Omega_1} d\Omega_1 d\tau, \quad (9)$$

where $N_0 \approx 1.37 \times 10^{15}$ is the number of particles in a bunch and $\Delta \tau = t / c \approx 2 \times 10^{-6}$ s is the flight time in the LEP straight section $l \approx 600$ m [9]. The three ICS spectra, derived from our equations and from the Monte Carlo simulation are shown in Fig. 2 and labeled respectively as:

- M — Monte Carlo simulation (Ref. [9], Fig. 6);
- T — Thomson approximation, from Eqn (8);
- C — Compton approximation, from Eqn (6).

![Figure 1](image1.png)

**Figure 1.** ICS spectrum from Eqn (8) for $T = 291$ K and non relativistic (BBR) $\gamma = 1$ (dot), and relativistic $\gamma = 10$ (dot dash), $\gamma = 10^2$ (dash), $\gamma = 10^3$ (continuous) Lorentz factor.

![Figure 2](image2.png)

**Figure 2.** ICS spectra at LEP: Monte Carlo simulation (M) (Ref. [9], Fig. 6), Thomson approximation (T) [Eqn (8)], Compton approximation (C) [Eqn (6)].
Note that our Compton spectrum shows good agreement with the Monte Carlo simulation while the Thomson one, at high energies, is a factor of 3 higher. This overestimate of the Thomson spectrum is clearly related to the independence of the Thomson cross section with respect to the photon energy. We notice that our ICS spectrum, in the Compton limit, is nearly coincident with the Monte Carlo approximation at low energy but for higher energies (> GeV) it is 26% above. This discrepancy seems not to be related to our approximations because, from our data we obtain, for the total event number, a value \( N_1 = 2.646 \) practically coincident (within 0.1%) with \( N_1 = 2.65 \) found by Di Domenico [9] by Monte Carlo simulations. The difference might be due to the Monte Carlo method whose statistical procedure implies a smaller number of events at higher energies. However this discrepancy does not affect the beam lifetime because only the number rate \( dN_1/dt_1 \) is involved in its evaluation. We recall that ICS may also be used to study the bunch internal structure and the result of coherent emission by the charges, in this case the optimal experimental set up is reached when the relativistic bunches are hit by a collinear back photon emitted by a laser [11].

4. ICS in Astrophysics

\( X - \gamma \) astronomy mostly traces the presence of relativistic electrons (or, at lower level, of relativistic nuclei) by their synchrotron radiation or their ICS off infrared, interstellar or cosmic radiation. Moreover the same ICS may become the main slowing process for relativistic charges once the magnetic field energy densities \( \rho_B \) are below the corresponding cosmic photon energy densities \( \rho_{BBR} \). This situation generally occurs in extragalactic space where \( \rho_B \ll \rho_{BBR} \). In this framework cosmic ray electrons around SN1006 have been recently discovered indirectly by their non-thermal X-ray imaging cosmic photon energy densities [12]. It is therefore of great interest and pertinence to provide not just an order of magnitude for the corresponding ICS \( \gamma \) ray flux but also for its detailed spectrum for such relativistic and ultrarelativistic cosmic rays electrons. We show in Fig. 3 these ICS spectra for \( \gamma = 10^4, 10^6, 10^8 \). It is important to note the existence of a high energy transition in the ICS from Thomson to Compton behavior. This change occurs dramatically at the highest energy range of the spectra. The Lorentz factor \( \gamma \) for Compton behavior occurs for electrons \( (e^- + \gamma \rightarrow e^- + e^+ + e^-) \) and for protons \( (\gamma + p \rightarrow n + \pi^+, p + \pi^0 \ldots) \) in ICS off \( T = 2.73 \text{ K} \) BBR respectively at the huge energy values

\[
\begin{align*}
\gamma_e &= \frac{2m_e c^2}{m_e} \frac{1}{\epsilon_0} \approx 1.6 \times 10^9, \\
\gamma_p &= \frac{2m_e c^2}{m_p} \frac{1}{\epsilon_0} \approx 2.4 \times 10^{11}.
\end{align*}
\]

ICS behavior in the Compton limit may be qualitatively predicted keeping in mind that the ICS energy spectrum falls off at photon energies near \( \epsilon_1 \approx \gamma^2 \kappa g T \) and the energy conservation calls for a cut-off of the extension of the ICS spectrum plateau from \( \kappa g T \) to \( mc^2 \) instead of reaching the usual extreme value \( \gamma^2 \kappa g T \). Let us better understand this ICS behavior change from Thomson to Compton regime as follows: as long as the incident photon, in the EF, has a characteristic energy \( \epsilon_0 \approx \gamma \kappa g T < mc^2 \) the diffused photon, in the EF, maintains most of its original energy and is spread nearly isotropically. Therefore, once the spectrum is reviewed in the LF after the Lorentz boost, the previous different angular distribution of the photons in the EF becomes, in the LF, a different energy distribution of the same photons. The angular integral of these differential ‘BBR’ spectra \( dN_1/(d\epsilon_1 \, d\epsilon_0 \, d\Omega) \), whose ‘Rayleigh’ regions overlap, produces the wider plateau shown in Fig. 4 in the energy range \( \kappa g T < \epsilon_1 < \gamma^2 \kappa g T \). However, in the Compton regime, when \( \gamma \kappa g T > mc^2 \) the differential cross section \( \sigma \sim \sigma_T \, m_1/\epsilon_0^2[1 - \cos \theta_{sc}] \) leads, in the EF, to a diffused and anisotropic photon distribution which becomes more and more beamed into a ‘Compton cone’ at angles \( \theta_{c} \ll \sqrt{me^2/(\gamma \kappa g T)} \sqrt{1/\gamma} \) (not to be confused with the thinner ‘Lorentz’ angles \( \theta_{L} \ll 1/\gamma \) related to the boost transformations). All the photons diffused outside the \( \theta_{c} \) cone can reach final energies of order \( \epsilon_1 \approx \gamma^2 \kappa g T(1 - \beta \cos \theta_{sc}) \approx 2 \gamma^2 \kappa g T(1 - \beta \cos \theta_{sc}) \ll mc^2/\gamma \) in the LF. On the other side, around and inside the Compton cone \( \theta_{c} \) but far from the edge of the inner Lorentz cone \( \theta_{L} \) (for example we may consider \( 1/2 \sqrt{me^2/(\gamma \kappa g T)} \ll \theta_{L} < \sqrt{me^2/(\gamma \kappa g T)} \) where the majority (>75%) of the scattered photons are found) the final available energy, in the LF, will be \( \epsilon_1 \approx \gamma^2 \kappa g T(1 - \beta \cos \theta_{sc}) \approx \kappa g T + (mc^2/8\gamma) \). In conclusion the Compton behavior of the ICS spectra, derived analytically from the exact formula (5) and described in Fig. 5, pile up the photons near the edge energies in the range \( \epsilon_1 \geq mc^2/\gamma \) and leads to an unexpected peak higher and higher at those highest energy values as the Lorentz factor \( \gamma \geq mc^2/\kappa g T \) increases. It is also important to notice that this extreme

![Figure 3. ICS spectrum for \( T = 2.73 \text{ K} \) and \( \gamma = 10^4 \) (dot dash), \( \gamma = 10^6 \) (dot), \( \gamma = 10^8 \) (continuous) Lorentz factor.](image)

![Figure 4. ICS ‘Thomson’ spectrum for \( T = 2.73 \text{ K} \) and a \( \gamma = 10^4 \) Lorentz factor.](image)
Compton regime may arise in a different way: for example in non-relativistic or relativistic cases when $\gamma_b T \ll mc^2$.

5. The main scenarios for ICS applications

For a synthetic but complete picture of all the above described ICS behavior we suggest defining the following characteristic regimes each labeled by the relevant parameters involved.

(1) ($\gamma \sim 1, \gamma_b T \ll mc^2/\gamma$) The non-relativistic Compton scattering limit off cold-warm BBR (Fig. 6). Here the simplest ICS spectrum becomes a self similar Planck spectrum; its reflective efficiency depends on the usual quantities $n_e, \sigma_T, ...$. One of the most celebrated applications is the Sunyaev–Zeldovich effect in cosmology.

(2) ($\gamma \gg 1, \gamma_b T < mc^2/\gamma$) ICS at the relativistic Thomson limit off cold-warm BBR (Figs 1, 2). This is the case with most application in high energy physics and astrophysics (LEP I, LEP II, GRBs, cosmic rays at energies $E_e \ll 10^{14}$ eV and $E_p \ll 10^{15}$ eV,...). Now the ICS spectrum deviates from a pure Planckian spectrum leading to a ‘cut-hill’ spectrum with a smooth edge covering the energies from $\epsilon_1 \sim \gamma_b T$ up to photon energies $\epsilon_1 \sim \gamma \gamma_b T$. This ‘simple’ ICS spectrum may rule the spectra of ICS by charge (electron) jets ($\gamma \sim 10^{-4} \sim 10^4$) onto nearby stellar companion photons ($\epsilon_0 \sim 0.5$ eV) leading to successful gamma jets able to explain the integral spectra of GRBs [13]. This argument has been widely developed by the authors and is still under consideration [14]. Because of the ‘smooth’ behavior of the edge of the spectrum there is not a one to one relation between the cosmic rays electron or proton spectral index and the final ICS $\gamma$ rays spectral index. Another scenario where ICS leads to $\gamma$ rays arises in extragalactic blazars which eject beamed cosmic rays (protons, nuclei,...) over huge distances. Their ICS off 2.73 K BBR may also lead to collinear $\gamma$ jets; their energy is dominated by the electron presence (over the nuclei one) in the cosmic rays jets. Such a gamma jet is the large scale version of the minijet model plus ICS considered by the authors for GRB production in galactic binary systems.

(3) ($\gamma \gg 1, mc^2/\gamma \ll \gamma_b T \ll mc^2$) ICS at the relativistic Compton limit off cold-warm BBR (Fig. 7). These energy windows are relevant for the highest cosmic ray interactions off BBR ($E_e \gg 10^{14}$ eV, $E_p \gg 10^{30}$ eV). The ICS spectrum exhibits a pile up of photons at the edge of the highest $\epsilon_1$ energies leading to a peaked maximum at $\epsilon_1 \sim mc^2\gamma$ and a sharp cut off at energies just above it. The presence of such a peak has important consequences in the cosmic rays $\gamma$ rays links. The primordial incident cosmic ray spectrum at highest energies (power law, ...) leaves its original imprint on a similar ‘photocopy’ $\gamma$ ray spectrum. Moreover the new born high energy $\gamma$ rays ($\epsilon_1 > 100$ TeV) may also successfully interact with BBR photons (electron pair production) leading to electron-proton cascades or electromagnetic showers at lower and lower energies. The showers will arrest their growth as soon as the last degraded photon energies reach the threshold $E_e = [2mc^2/\gamma_b T]mc^2 = 10^{15}$ eV. It is important to recall that the $\gamma - \gamma \rightarrow e^+ + e^-$ process at the ultrarelativistic limit has a very similar cross section (Breit–Wheeler, 1934) to $e^- e^-$ annihilation (Dirac, 1930) or the Klein–Nishina (1929) cross section. One of the consequences is the expected cosmic background presence of relic TeV $\gamma$ rays components (due to the direct cosmological ICS or due to secondary showers at energies $\epsilon_1 \lesssim 100$ TeV). This argument will be discussed in detail elsewhere.

(4) ($\gamma \sim 1, \gamma_b T \gg mc^2$) Non-relativistic Compton scattering off ‘hot’ BBR (Fig. 8). The resulting ICS spectrum deviates from the original and exhibits a small peak at energies $\epsilon_1 \sim mc^2/2$. The effect as above is related to the energy dependence of the Klein–Nishina cross section.

(5) ($\gamma \gg 1, mc^2 \ll \gamma_b T \ll mc^2$) Relativistic ICS off ‘hot’ BBR. The spectrum shows a new ‘plateau’ extending from $\epsilon_1 \sim \gamma_b T$ energies up to $\epsilon_1 \sim mc^2\gamma$ energies with a marked peak at $\epsilon_1 \sim mc^2\gamma$ energies and a sharp cut off above (this behavior is similar to point 3).
The equilibrium regime of ICS off ‘hot’ BBR (Fig. 9). The spectrum shows a marked peak at $\epsilon_1 \sim mc^2 \gamma$ energies with a sharp cut-off above. The known astrophysical scenarios where such ICS may play a role could be the earliest hottest (thermal equilibrium) cosmological epochs ($t \lesssim 1 \ s$, $T > \text{MeV}$) and in the hot thermal cores of supernovae explosions ($x_B T \gtrsim 5$). It is interesting to notice the nature of the non equilibrium ICS spectrum and it may be worthwhile deriving the exact kinetic equations (due to such ICS) for the multicomponent thermal bath of the early universe as well as the SN explosion process. Finally if the fireball model is a real event as needed to explain the GRB puzzle then such ICS (off the last layers of the fireball explosion) would be smeared by multiscattering during last stages of the fireball into a final near-thermal GRB spectrum. We do not recognize the presence of such a thermal imprint in GRB [13].

6. Conclusions and applications

The ICS formulae developed have a wide range of applications; in particular in fitting LEP I, LEP II experimental data and in understanding recent puzzling GRB spectra (in this last case we modified Eqn (8) in order to take into account the diluted and anisotropic BBR spectrum seen from the jet and we assumed a ring-like photon source [Ref. [14], Eqn (20)]. Moreover the extreme spectra of ICS at the relativistic Compton limit off a cold-warm BBR, its peak at very high energies, may be probed at LEP I, LEP II using a diffused thermal optical light (a flash) in the beam pipe during the bunch crossing. Finally the recent evidence for cosmic ray electrons at energies above hundreds TeV by their observed synchrotron radiation in soft X spectra implies the coexistence of a low but nearly detectable component of high $\gamma$ cosmic rays at energies $E_{\gamma} \lesssim mc^2 \gamma \approx 100 \ \text{TeV}$. Because of the arguments in Section 3 the cosmic ray electron spectrum ($dN/dE_{\gamma} \sim E_{\gamma}^{-2}$) will also be reflected in the ‘photocopy’ $\gamma$ spectrum. The total energy flux $\phi_\gamma$ will be of the same order of magnitude as the X rays flux:

$$\phi_\gamma \lesssim \phi_X \frac{\rho_{\text{BBR}}}{\rho_B} \lesssim \frac{1}{2} \frac{\phi_X}{\epsilon_{\gamma}} \left( \frac{B}{6 \ \mu \text{G}} \right)^{-2}. \quad (11)$$

The corresponding $\gamma$ ray flux number will be extremely suppressed because of the $X - \gamma$ energy ratio,

$$\frac{dN_\gamma}{dN_X} \sim \frac{\phi_\gamma}{\phi_X} \left( \frac{E_{\gamma}}{E_X} \right) \sim 2 \times 10^{-11} \left( \frac{B}{6 \ \mu \text{G}} \right)^{-1} \left( \frac{E_{\gamma}}{100 \ \text{TeV}} \right). \quad (12)$$

and it follows that the area needed to detect such a low flux is possibly below the present air shower arrays sensitivity. However an accurate direction–flux correlation might be able in the near future to observe a tiny 100 TeV $\gamma$ flux. We suggest seriously considering, a detailed observational program in order to detect the above tiny 100 TeV $\gamma$ flux from SN1006 and possibly the other suggested candidates (Cas A, IC443, Tycho SN, ...) of SNRs. Their presence in the above fluxes is a necessary and compelling consequence of fundamental QED and cosmological BBR theories combined in present ICS models. Finally we interpret the two X lobes around SN1006 as being generated by an electron jet...
scattering off the relic giant shell contrary to the more popular idea of a Fermi shock acceleration mechanism. In these beaming models one should also expect a rare and strongly time dependent TeV $\gamma$ ray 'burst' which could be observed over a short period of time (hours) once the observer is inside the thin jet cone direction, with an integral intensity amplified by Lorentz factor $\gamma$ over a diffused spherical source. Well known candidates are blazars or quasars such as 3C279, AGNs, such as NGC 3079, or most recent TeV Mrk sources. Similar arguments lead us to expect a low variable relic background of gamma (tens of TeV) noise due to the pile up of cosmic integral ICS gamma rays and their electromagnetic cascades just below the electron pair creation threshold off cosmic BBR.

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