

ASTROPHYSICAL GAMMA-RAY PRODUCTION BY INVERSE COMPTON INTERACTIONS OF RELATIVISTIC ELECTRONS

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ABSTRACT

The inverse Compton scattering of background photon gases by relativistic electrons is a good candidate for the production of high-energy γ -rays in the diffuse interstellar medium as well as in discrete sources. By discussing the special case of the scattering of the diffuse starlight in the interstellar medium by cosmic ray electrons, we demonstrate that previous derivations of the γ -ray source function for this process on the basis of the Thomson limit of the Klein-Nishina cross section lead to incorrect values for γ -ray energies above 100 MeV. It is shown that the Thomson limit is not applicable for the calculation of γ -ray source functions in astrophysical circumstances in which target photons with energies greater than 1 eV are scattered by relativistic electrons.

Subject headings: atomic processes — gamma rays: general — interstellar: matter

I. INTRODUCTION

The recent spectral measurements of the galactic γ -radiation with energies between 50 MeV and 4 GeV on board the European satellite experiment *COS B* (Paul *et al.* 1978) as well as balloon-borne observations of the medium γ -ray range (Kniffen *et al.* 1978) have again raised the topic of the origin of this radiation since the measured spectra do not resemble a pure pion spectrum which has been predicted theoretically (Dodds *et al.* 1975). The measured, rather steep spectra suggest more an electromagnetic origin of the galactic γ -radiation induced by cosmic ray electrons. These electromagnetic processes may take place in discrete point sources as well as in interstellar space.

A process which is a good candidate for the production of high-energy γ -rays in the diffuse interstellar medium as well as in discrete sources is the inverse Compton scattering of background photon gases by relativistic electrons. In the interstellar medium, mainly photons of the universal 2.7 K microwave background radiation, the infrared radiation emitted by galactic dust and the diffuse starlight are scattered by the electrons into the γ -ray regime. In the context of discrete sources one may think of the Compton-synchrotron effect which has already been applied to the X-ray and γ -ray emission from pulsars (Apparao and Hoffman 1970; Grindlay and Hoffman 1971; Zheleznyakov and Shaposhnikov 1972), compact radio sources (Rees 1967; Albano, Terlevich, and Frank 1973), and Seyfert galaxies (Bergeron and Salpeter 1971, 1973).

In order to estimate quantitatively the importance of the inverse Compton scattering as a γ -ray production process a detailed knowledge of the γ -ray source function is needed on the basis of a given electron distribution and a given target photon distribution. By discussing the special case of the scattering of the diffuse starlight in the interstellar medium we demonstrate that previous derivations of this γ -ray source function on the basis of the Thomson limit of the Klein-Nishina cross section lead to incorrect values for γ -ray energies above 100 MeV. It is shown that the Thomson limit is not applicable in astrophysical circumstances in which target photons with energies greater than 1 eV are scattered by relativistic electrons.

II. INVERSE COMPTON EMISSIVITY

The kinematics of the inverse Compton effect has extensively been discussed in the literature (Jones 1968; Blumenthal and Gould 1970). In this section we follow the notation of Blumenthal and Gould (BG). We assume that the distribution of the target photons $n(\epsilon, \mathbf{r})$ at a galactic position $\mathbf{r}$ is spatially isotropic and that the incident electrons have relativistic energies (Lorentz factor $\gamma = E/mc^2 \gg 1$). Then the omnidirectional differential source function of inverse Compton scattered target photons is given by

$$S(E_\gamma, \mathbf{r}) = 8\pi^2 r_0^2 (mc^2)^2 \int_{E_{\min}}^{\infty} dE E^{-2} I(E, \mathbf{r}) \int_0^{\infty} d\epsilon \epsilon^{-1} n(\epsilon, \mathbf{r}) \\ \times \left[2q \ln q + (1 + 2q)(1 - q) + \frac{1}{2} \frac{(\Gamma_e q)^2 (1 - q)}{(1 + \Gamma_e q)} \right] \text{GeV}^{-1} \text{cm}^{-3} \text{s}^{-1}, \quad (1)$$

where

$$\Gamma_e = 4\epsilon\gamma/mc^2, \quad (2)$$

$$q = E_1/[\Gamma_e(1 - E_1)], \quad (3)$$

$$E_1 = E_\gamma/\gamma mc^2, \quad (4)$$

r_0 is the classical electron radius, and $I(E, \mathbf{r})$ denotes the differential cosmic ray electron intensity at the position $\mathbf{r}$. In astrophysics the inverse power law is one of the most commonly occurring electron energy distributions (Jones 1968), so we shall consider the electron spectrum

$$\begin{aligned} I(E, \mathbf{r}) &= N(\mathbf{r})E^{-p} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} & \text{for } E \geq E_c \\ &= 0 \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} & \text{for } E < E_c, \end{aligned} \quad (5)$$

where the cutoff energy E_c is considerably larger than mc^2 , say $E_c = 10mc^2$, and the spectral index p is a positive constant. This form of the electron intensity distribution is suggested by most of the proposed cosmic ray acceleration mechanisms and is confirmed to hold over wide electron energy ranges by studies of the frequential distribution of nonthermal radio emission from probable cosmic ray sources like pulsars and supernova remnants and the galactic background (Alexander and Clark 1974) as well as observations of the local demodulated cosmic ray electron spectrum (Ramaty 1974).

Assuming that we are away from the endpoint of the Compton spectrum ($E_\gamma > E_c$) so that we can neglect "cutoff effects" in the integration over electron and initial photon energies, expression (1) reduced to (BG)

$$S(E_\gamma, \mathbf{r}) = 2^{p+3}\pi^2 N(\mathbf{r}) r_0^2 (mc^2)^{1-p} E_\gamma^{-(p+1)/2} \int_0^\infty d\epsilon n(\epsilon, \mathbf{r}) \epsilon^{(p-1)/2} B[s(\epsilon, E_\gamma), p] \text{ GeV}^{-1} \text{ cm}^{-3} \text{ s}^{-1} \quad (6)$$

with the integral

$$B[s(\epsilon, E_\gamma), p] = \int_0^1 dq q^{(p-1)/2} \frac{2q \ln q + 1 + q - 2q^2 + 2sq(1-q)}{\{1 + [sq/(1+sq)]^{1/2}\}^{p+2} (1+sq)^{(p+3)/2}} \quad (7)$$

and the dimensionless parameter

$$s = \epsilon E_\gamma / m^2 c^4 \quad (8)$$

which determines the domain of scattering.

For values of s smaller than 1, the integrand in (7) can be expanded by keeping only the lowest order terms in $(qs)^{1/2}$, so that the q -integration can be performed. The first four order terms are given in the Appendix. With the zeroth order term $B^{(0)}(p)$ one derives the well-known approximation of (6) in the Thomson limit.

III. SCATTERING OF STARLIGHT PHOTONS: VALIDITY OF THE THOMSON LIMIT

We now apply the general results of § II to the scattering of the local diffuse starlight photons by cosmic ray electrons. Since we restrict ourselves to astrophysical circumstances in the local spiral arm, we may neglect possible large-scale $\mathbf{r}$ -dependences of the starlight energy density and the cosmic ray electron density. It has been pointed out by Stecker (1977) that the local starlight distribution can be represented by the sum of two gray-body components of equal energy density $w_I = w_{II} = 0.22 \text{ eV cm}^{-3}$. The first part with a temperature of $T_I = 10^4 \text{ K}$ will hereafter be referred to as Population I stars; the second part with a temperature $T_{II} = 5.10^3 \text{ K}$ is called the Population II component.

With these target photon distributions,

$$n_{I,II}^L(\epsilon) = \frac{15w_{I,II}}{\pi^4(kT_{I,II})^4} \frac{\epsilon^2}{\exp(\epsilon/kT_{I,II}) - 1}, \quad (9)$$

the local differential omnidirectional source function can be calculated from (6) either by analytical integration using the expansions (A1)–(A4) or by direct numerical computation. The former method gives

$$S_L(E_\gamma) = S_I(E_\gamma) + S_{II}(E_\gamma) \quad (10)$$

with

$$S_{I,II}(E_\gamma) = \frac{45\sigma_T w_{I,II}}{2\pi^3} N_L(mc^2)^{1-p} F(p)(kT_{I,II})^{(p-3)/2} E_\gamma^{-(p+1)/2} f_{\text{corr}}(p, E_\gamma, T_{I,II}) \text{ GeV}^{-1} \text{ s}^{-1} \text{ cm}^{-3}, \quad (11)$$

where $\sigma_T = 6.65 \times 10^{-25} \text{ cm}^2$ is the Thomson cross section. The function

$$F(p) = 2^{p+3} \Gamma\left(\frac{p+5}{2}\right) \zeta\left(\frac{p+5}{2}\right) \frac{p^2 + 4p + 11}{(p+1)(p+5)(p+3)^2} \quad (12)$$

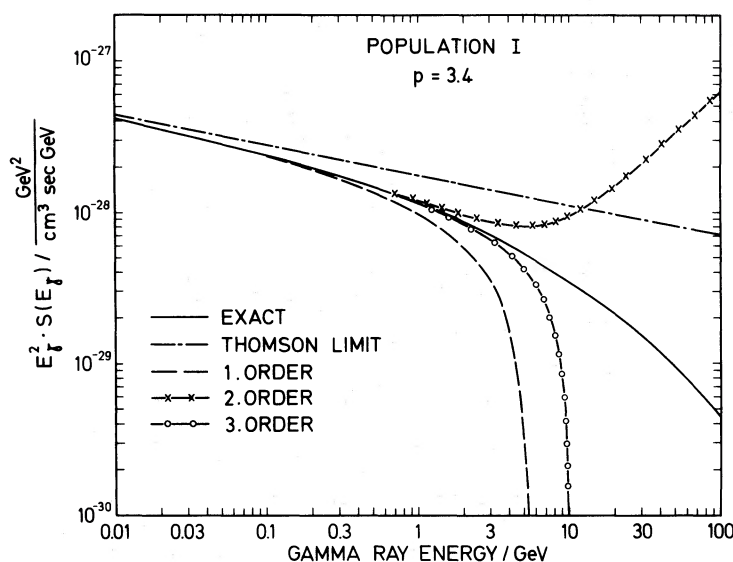


FIG. 1.—Omnidirectional differential γ -ray source function multiplied by E_γ^2 for Population I photons calculated in several approximations in comparison with the exact result.

is tabulated by BG. $\Gamma(x)$ and $\zeta(x)$ denote the gamma function and Riemann's zeta function, respectively. The function $f_{\text{corr}}(p, E_\gamma, T_{\text{I,II}})$, which is equal to unity in the Thomson limit, contains the corrections to the Thomson limit.

With the first-order expansion (A2) BG obtained

$$f_{\text{corr}}^{(1)}(p, E_\gamma, T) = 1 - \left\{ \frac{(p^2 + 6p + 16)(p + 1)(p + 5)(p + 3)^2}{(p^2 + 4p + 11)(p + 6)(p + 4)^2} \frac{\Gamma[\frac{1}{2}(p + 6)]\zeta[\frac{1}{2}(p + 6)]}{\Gamma[\frac{1}{2}(p + 5)]\zeta[\frac{1}{2}(p + 5)]} \frac{(E_\gamma kT)^{1/2}}{mc^2} \right\}. \quad (13)$$

The higher-order correction functions $f_{\text{corr}}^{(2)}(p, E_\gamma, T)$ and $f_{\text{corr}}^{(3)}(p, E_\gamma, T)$ are given in the Appendix.

It is interesting to note that these correction functions do not depend on the target photon energy density w since w enters equation (9) only as a multiplicative factor independent of ϵ .

Now, let us discuss the results concerning the omnidirectional differential and integral source functions. In Figures 1 and 2 the differential source functions for both populations are shown calculated according to equation

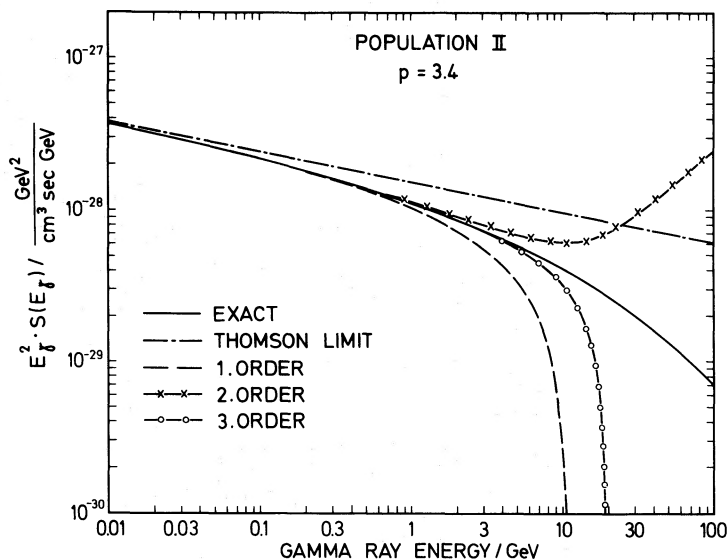


FIG. 2.—Omnidirectional differential γ -ray source function multiplied by E_γ^2 for Population II photons calculated in several approximations in comparison with the exact result.

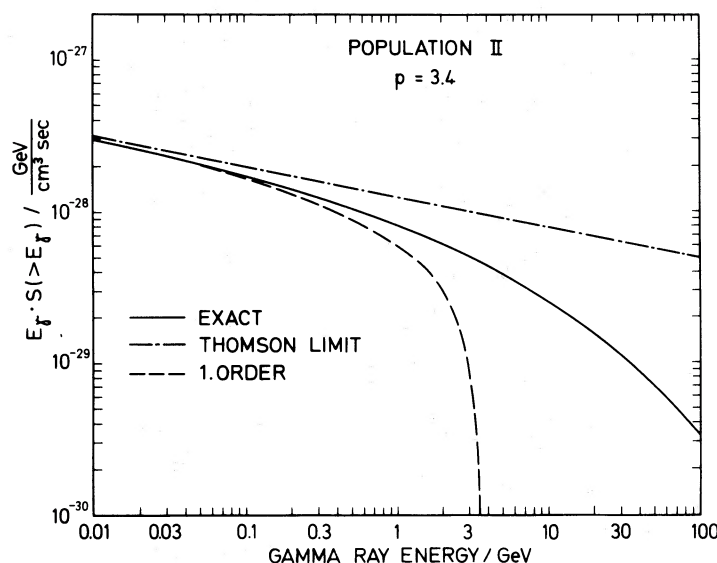


FIG. 3.—Omnidirectional integral γ -ray source function multiplied by E_γ for Population II photons calculated in the first two approximations in comparison with the exact result.

(11) as well as by direct numerical integration of equation (6). In our calculations we use the cosmic ray electron spectrum parameters $p = 3.4$ and $N_L = 0.08 \text{ GeV}^{2.4} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$ measured by Meegan and Earl (1977).

In both cases it can be seen that neither the Thomson limit nor the lowest-order calculations are able to reproduce the exact values of the differential source function over the whole γ -ray energy range. The Thomson limit gives exact results only for γ -ray energies lower than 100 MeV; the first-order calculations can be used up to energies of 1 GeV whereas the third-order results are restricted to γ -ray energies below 5 GeV (Population I) and 10 GeV (Population II).

Therefore, the results of previous calculations of the differential Compton source function which were based entirely on the Thomson limit (Ginzburg and Syrovatskii 1964; Felten and Morrison 1966; Stecker 1971; Dodds *et al.* 1975; Shukla and Paul 1976; Piccinotti and Bignami 1976) are valid only for the lower-energy γ -ray regime ($E_\gamma < 100 \text{ MeV}$). In the range of $100 \text{ MeV} < E_\gamma < 5 \text{ GeV}$ one may use the expansions in the near Klein-Nishina limit.

Concerning the integral source functions these conclusions become even stronger. In Figure 3 we show the omnidirectional integral Compton source function for Population II photons calculated from equation (11) to zeroth and first order as well as by direct numerical integration of (6). It can be seen that both the Thomson limit and the first-order calculation lead to large deviations of the integral source function from the exact one, particularly at higher energies $E_\gamma > 500 \text{ MeV}$. But even at 100 MeV the Thomson limit value is a factor of 1.17 too high.

To make clear that this inaccuracy of the Thomson limit holds also for Population I photons or other representations of the local starlight distribution, we show in Table 1 that γ -ray energy E_γ^0 at which the first-order correction function (13) is equal to 0.5 at fixed p and fixed T . This γ -ray energy is a reasonable upper bound for using the lowest-order expansion results. We find that for all possible representations of the starlight spectrum and for all possible values of the electron spectral index p neither the Thomson limit nor the lowest-order expansion calculations are applicable to the whole γ -ray range. For the case of high-energy γ -ray source functions, it therefore seems to be more appropriate to integrate expression (6) numerically.

Furthermore, it is interesting to note that the exact differential and integral Compton γ -ray spectra are no longer pure power laws as in the Thomson limit. Concerning the differential spectrum we have a continuous steepening from the Thomson behavior at lower γ -ray energies (spectrum $\propto E_\gamma^{-(p+1)/2}$) the onset of which depends mainly upon the mean energy of the target photon gas $\langle \epsilon \rangle = 2.7kT$.

IV. DISCUSSION OF THE COMPTON SCATTERING OF GENERAL PHOTON GASES

In order to generalize our conclusions to the case of the scattering of other photon gases in the interstellar space, one must specify the spectrum of these target photons, so that the ϵ -integration in (6) can be carried out. However, in most astrophysical circumstances these spectra are only poorly known, so that an exact analysis as in § III is not possible. But, on the other hand, some qualitative conclusions can be drawn.

Our discussion of the starlight scattering in § III has shown that the electron spectral index p and the mean energy of the starlight photons ($2.7kT$) are the basic parameters which determine the behavior of the Compton

TABLE 1
UPPER-BOUND GAMMA-RAY ENERGY (in GeV) FOR USING THE LOWEST-ORDER EXPANSION RESULTS FOR FIXED p AND T

p	$T[\text{K}]$											
	2×10^3	4×10^3	6×10^3	8×10^3	1×10^4	1.2×10^4	1.5×10^4	1.8×10^4	2×10^4	3×10^4	4×10^4	5×10^4
1.4.....	37.2	18.6	12.4	9.3	7.4	6.2	5.0	4.1	3.7	2.5	1.9	1.5
1.6.....	29.5	14.8	9.8	7.4	5.9	4.9	3.9	3.3	3.0	2.0	1.5	1.2
1.8.....	23.9	11.9	8.0	6.0	4.8	4.0	3.2	2.7	2.4	1.6	1.2	0.95
2.0.....	19.6	9.8	6.5	4.9	3.9	3.3	2.6	2.2	2.0	1.3	0.98	0.78
2.2.....	16.3	8.1	5.4	4.1	3.3	2.7	2.2	1.8	1.5	1.4	0.81	0.65
2.4.....	13.7	6.8	4.6	3.4	2.7	2.3	1.8	1.5	1.4	0.91	0.68	0.55
2.6.....	11.6	5.8	3.9	2.9	2.3	1.9	1.6	1.3	1.2	0.78	0.58	0.47
2.8.....	10.0	5.0	3.3	2.5	2.0	1.7	1.3	1.1	1.0	0.66	0.50	0.40
3.0.....	8.6	4.3	2.9	2.2	1.7	1.4	1.1	0.96	0.86	0.57	0.43	0.34
3.2.....	7.5	3.7	2.5	1.9	1.5	1.2	1.0	0.83	0.75	0.50	0.37	0.30
3.4.....	6.6	3.3	2.2	1.6	1.3	1.1	0.88	0.73	0.66	0.44	0.23	0.26
3.6.....	5.8	2.9	1.9	1.4	1.2	0.97	0.77	0.64	0.58	0.39	0.29	0.23
3.8.....	5.1	2.6	1.7	1.3	1.0	0.86	0.68	0.57	0.51	0.34	0.26	0.21
4.0.....	4.6	2.3	1.5	1.1	0.91	0.76	0.61	0.51	0.46	0.30	0.23	0.18

source functions in the non-Thomson domain. Therefore, it seems reasonable to work with these quantities for the case of scattering photon gases with poorly known spectra. In practice, we approximate the real unknown target photon distribution by an idealized monochromatic distribution

$$n(\epsilon) = \frac{w_{ph} \delta(\epsilon - \langle \epsilon \rangle)}{\langle \epsilon \rangle} \text{ photons GeV}^{-1} \text{ cm}^{-3}, \quad (14)$$

where

$$w_{ph} = \int_0^\infty d\epsilon \epsilon n(\epsilon) \text{ GeV cm}^{-3} \quad (15)$$

is the photon energy density and

$$\langle \epsilon \rangle = w_{ph} / \int_0^\infty d\epsilon n(\epsilon) \text{ GeV} \quad (16)$$

denotes the mean energy.

With these approximations it is straightforward to repeat the analysis of § III: e.g., for the first-order correction function we find now

$$f_{\text{corr}}^{(1)}(p, E_\gamma, \langle \epsilon \rangle) = 1 - \frac{(p^2 + 6p + 16)(p + 1)(p + 5)(p + 3)^2}{(p^2 + 4p + 11)(p + 6)(p + 4)^2} \left(\frac{E_\gamma \langle \epsilon \rangle}{m^2 c^4} \right)^{1/2}. \quad (17)$$

Analogous to Table 1 we show in Figure 4 that γ -ray energy E_γ^0 at which this first-order correction function (17) is equal to 0.5 for different values of p as a function of the mean photon energy $\langle \epsilon \rangle$ covering the frequency range from low-frequency radio photons to soft X-ray photons. It can be seen that for all values of p —in order to calculate γ -ray source functions from inverse Compton scattering at $E_\gamma \geq 100$ MeV—the Thomson limit of the Klein-Nishina cross section is not applicable if target photon gases with mean energies greater than 1 eV are involved.

On the other hand, it is clear from our discussion in § III that the appearance of a continuous steepening from the power law behavior of an observed γ -ray energy spectrum favors an inverse Compton origin of the γ -rays in which optical, ultraviolet, or X-ray target photons—depending on the onset energy of the steepening—are

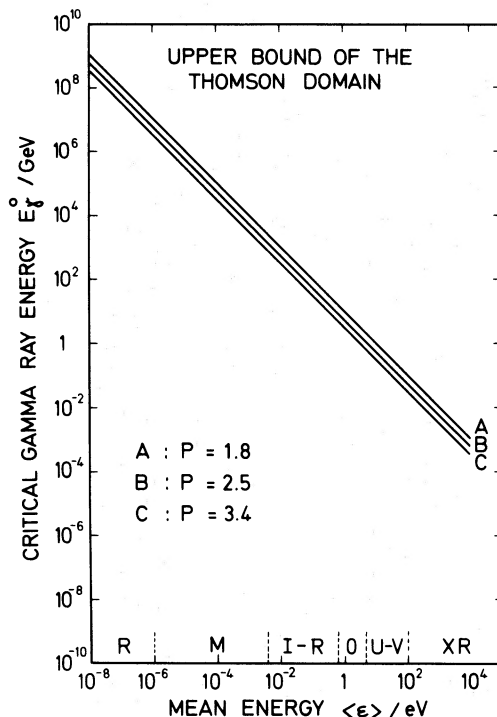


FIG. 4.—Critical γ -ray energy E_γ^0 which is a reasonable upper bound for using the Thomson limit approximation as a function of the mean energy $\langle \epsilon \rangle$ of the target photon distribution for different values of the electron spectral index p . The values of $\langle \epsilon \rangle$ cover the spectral range from radio (R), microwave (M), infrared (I-R), optical (O), ultraviolet (U-V), and X-ray (XR) wavelengths.

scattered into the γ -ray regime. In this context, it is interesting to note that such γ -ray spectra have been observed. Schönfelder's (1978) measurements of the γ -ray spectrum of the Seyfert galaxy NGC 4151 show strong evidence for a steepening of the spectral slope around 3 MeV. Secondly, the *COS B* observations of the pulsed γ -ray emission from the Vela pulsar PSR 0833-45 (Bennett *et al.* 1977) also indicates a steepening from the power law behavior of the γ -ray energy spectrum around 600 MeV. The inverse Compton origin of these γ -rays would provide a quite natural explanation of the observed breaks resulting from higher order effects in the Klein-Nishina cross section. In a subsequent work, I will discuss the γ -ray production in these sources in more detail.

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APPENDIX: CORRECTIONS TO THE THOMSON LIMIT

Expanding the integrand in (7) by keeping only the lowest-order terms in $(qs)^{1/2}$, after performing the q -integration, one obtains

$$B^{(0)}(s, p) = \frac{4(p^2 + 4p + 11)}{(p + 1)(p + 5)(p + 3)^2}, \quad (\text{A1})$$

$$B^{(1)}(s, p) = B^{(0)}(s, p) \left[1 - \frac{(p^2 + 6p + 16)(p + 1)(p + 3)^2(p + 5)}{(p^2 + 4p + 11)(p + 4)^2(p + 6)} s^{1/2} \right], \quad (\text{A2})$$

$$B^{(2)}(s, p) = B^{(0)}(s, p) \left[1 - \frac{(p^2 + 6p + 16)(p + 1)(p + 3)^2(p + 5)}{(p^2 + 4p + 11)(p + 4)^2(p + 6)} s^{1/2} + \frac{(p^4 + 12p^3 + 62p^2 + 164p + 209)(p + 1)(p + 3)}{2(p^2 + 4p + 11)(p + 7)(p + 5)} s \right] \quad (\text{A3})$$

$$B^{(3)}(s, p) = B^{(0)}(s, p) \left[1 - \frac{(p^2 + 6p + 16)(p + 1)(p + 3)^2(p + 5)}{(p^2 + 4p + 11)(p + 4)^2(p + 6)} s^{1/2} + \frac{(p^4 + 12p^3 + 62p^2 + 164p + 209)(p + 1)(p + 3)}{2(p^2 + 4p + 11)(p + 7)(p + 5)} s - \frac{(p^4 + 14p^3 + 84p^2 + 296p + 576)(p + 1)(p + 2)(p + 5)(p + 3)^2}{6(p^2 + 4p + 11)(p + 4)(p + 8)(p + 6)^2} s^{3/2} \right]. \quad (\text{A4})$$

These results are valid for values of s smaller than 1. BG have already derived equations (A1) and (A2). For large values $s \gg 1$, BG derived

$$B^{(0)}(s \gg 1, p) \rightarrow 2^{-(p+1)} s^{-(p+1)/2} [\ln(s) + C(p)], \quad (\text{A5})$$

where $C(p)$ is a parameter of order unity. In the intermediate region $s \sim 1$, $B(s, p)$ has to be computed numerically.

With the gray-body target photon distribution (9) the ϵ -integration in (6) can be carried out by using the expansions (A1)–(A4). The effect of these higher order terms in s can be well described by the correction function $f_{\text{corr}}(p, E_\gamma, T)$ which is equal to unity in the Thomson limit. According to the expansions (A2)–(A4) we find

$$f_{\text{corr}}^{(1)}(p, E_\gamma, T) = 1 - \left\{ \frac{(p^2 + 6p + 16)(p + 1)(p + 5)(p + 3)^2}{(p^2 + 4p + 11)(p + 6)(p + 4)^2} \frac{\Gamma[\frac{1}{2}(p + 6)]\zeta[\frac{1}{2}(p + 6)]}{\Gamma[\frac{1}{2}(p + 5)]\zeta[\frac{1}{2}(p + 5)]} \frac{(E_\gamma kT)^{1/2}}{mc^2} \right\}, \quad (\text{A6})$$

$$f_{\text{corr}}^{(2)}(p, E_\gamma, T) = f_{\text{corr}}^{(1)}(p, E_\gamma, T) + \left\{ \frac{\Gamma[\frac{1}{2}(p + 7)]\zeta[\frac{1}{2}(p + 7)]}{\Gamma[\frac{1}{2}(p + 5)]\zeta[\frac{1}{2}(p + 5)]} \frac{(p^4 + 12p^3 + 62p^2 + 164p + 209)(p + 1)(p + 3)}{2(p^2 + 4p + 11)(p + 7)(p + 5)} \frac{E_\gamma kT}{(mc^2)^2} \right\}, \quad (\text{A7})$$

and

$$f_{\text{corr}}^{(3)}(p, E_\gamma, T) = f_{\text{corr}}^{(2)}(p, E_\gamma, T) - \left\{ \frac{\Gamma[\frac{1}{2}(p + 8)]\zeta[\frac{1}{2}(p + 8)]}{\Gamma[\frac{1}{2}(p + 5)]\zeta[\frac{1}{2}(p + 5)]} \frac{(p^4 + 14p^3 + 84p^2 + 296p + 576)(p + 1)(p + 2)(p + 5)(p + 3)^2}{6(p^2 + 4p + 11)(p + 4)(p + 8)(p + 6)^2} \frac{(E_\gamma kT)^{3/2}}{(mc^2)^3} \right\}. \quad (\text{A8})$$

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