The connection between thermal and non-thermal emission in gamma-ray bursts: general considerations and GRB 090902B as a case study

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ABSTRACT
Photospheric (thermal) emission is inherent to the gamma-ray burst (GRB) ‘fireball’ model. We show here that inclusion of this component in the analysis of the GRB prompt emission phase naturally explains some of the prompt GRB spectra seen by the Fermi satellite over its entire energy band. The sub-MeV peak is explained as multicolour blackbody emission, and the high-energy tail, extending up to the GeV band, results from roughly similar contributions of synchrotron emission, synchrotron self-Compton and Comptonization of the thermal photons by energetic electrons originating after dissipation of the kinetic energy above the photosphere. We show how this analysis method results in a complete, self-consistent picture of the physical conditions at both emission sites of the thermal and non-thermal radiation. We study the connection between the thermal and non-thermal parts of the spectrum, and show how the values of the free model parameters are deduced from the data. We demonstrate our analysis method on GRB 090902B: we deduce a Lorentz factor in the range $9 \leq \eta \leq 10^{7}$, photospheric radius $r_{\text{ph}} \simeq 7.2-8.4 \times 10^{11}$ cm and dissipation radius $r_{\gamma} \geq 3.5-4.1 \times 10^{15}$ cm. By comparison to afterglow data, we deduce that a large fraction $\epsilon_{\text{d}} \approx 85-95$ per cent of the kinetic energy is dissipated, and that a large fraction, $\sim$ equi-partition of this energy, is carried by the electrons and the magnetic field. This high value of $\epsilon_{\text{d}}$ questions the ‘internal shock’ scenario as the main energy dissipation mechanism for this GRB.

Key words: plasmas – radiation mechanisms: thermal – radiative transfer – scattering – gamma-ray burst: general.

1 INTRODUCTION
Although extensively studied for nearly two decades, the origin of the prompt emission of gamma-ray bursts (GRBs) is still puzzling. Up until these days, GRB prompt emission spectra are often modelled as a smoothly broken power law, which is referred to as the ‘Band’ function (Band et al. 1993; Preece et al. 1998a, 2000; Kaneko et al. 2006, 2008; Abdo et al. 2009b). While the ‘Band’ function often provides very good fits to the observed spectra over a limited energy range, it suffers two crucial drawbacks. In several bursts seen by the Fermi satellite, a high-energy tail extending up to tens of GeV was seen (e.g. GRB 090510, Ackermann et al. 2010; or GRB 090902B, Abdo et al. 2009a). The first drawback is that this tail requires more than the ‘Band’ function on its own to have an acceptable fit. However, the most severe drawback is that the ‘Band’ function, being empirical in nature, does not provide any information about the emission mechanism that produces the prompt radiation.

A common interpretation of the observed GRB spectrum is that it results from synchrotron emission, accompanied perhaps by synchrotron self-Compton (SSC) emission at high energies (Rees & Mészáros 1994; Sari & Piran 1997a; Pilia & Loeb 1998; Guetta & Granot 2003; Pe’er & Waxman 2004; Gupta & Zhang 2007). The emission follows the dissipation of a kinetic energy. The prevalent dissipation models involve either internal shocks (Rees & Mészáros 1994; Sari & Piran 1997b), magnetic energy dissipation in...
Poynting-dominated outflows (Thompson 1994; Spruit, Daigne & Drenkhahn 2001; Giannios & Spruit 2005; Zhang & Yan 2011), or collisional heating (Beloborodov 2010). These models have two main advantages. First, they can explain the complex light curves often seen (albeit with very little predictive power). Secondly, they naturally account for the non-thermal spectrum.

In spite of these advantages, in recent years evidence is accumulating for serious difficulties in these models. A well-known deficiency of the internal shock scenario is the low efficiency of energy conversion (Mochkovitch, Mattia & Marques 1995; Kobayashi, Piran & Sari 1997; Daigne & Mochkovitch 1998; Lazzati, Ghisellini & Celotti 1999; Guetta, Spada & Waxman 2001; Maxham & Zhang 2009), resulting from the fact that only the energy associated with the differential motion between the expanding ejecta shells after the GRB explosion can be dissipated, and that this energy is much lower than the energy associated with the bulk motion. This is in contrast to the observations which show high efficiency in \( \gamma \)-ray emission, of the order of tens of per cent (Zhang et al. 2007; Nysewander, Fruchter & Pe’er 2009). A second drawback is that optically thin synchrotron and SSC emission cannot account for the steepness of the low-energy spectral slopes (Criders et al. 1997; Preece et al. 1998b, 2002; Ghirlanda et al. 2003) (although part of the observed steepening may be accounted for when SSC in the Klein–Nishina limit is considered; see Daigne, Bosnjak & Dubus 2009; Bosnjak, Daigne & Dubus 2009). Thirdly, the high-energy spectral slope varies significantly from burst to burst (Preece et al. 1998a; Kaneko et al. 2006), with some bursts showing very steep high energy spectral slopes. This is in contradiction to the expected spectral slope of synchrotron emission from a power-law distribution of electrons, which is expected to produce a nearly flat spectrum, \( F(\nu) \propto \nu^0 \). Finally, the fitting is often made to the time-integrated spectrum. Analysis of time resolved spectra done by Criders, Liang & Preece (1998) and Ghirlanda et al. (2003) has shown that neither the synchrotron nor the SSC models can explain the time resolved low-energy spectral slopes.

These well-known difficulties of the synchrotron emission model have motivated the study of alternative scenarios. These include, among others, quasi-thermal Comptonization (Ghisellini & Celotti 1999), Compton drag (Zdziarski, Svensson & Paczyński 1991; Shemi 1994; Lazzati et al. 2000), jetter radiation (Medvedev 2000), Compton scattering of synchrotron self-absorbed photons (Panaitescu & Mészáros 2000; Stern & Poutanen 2004) and synchrotron emission from a decaying magnetic field (Pe’er & Zhang 2006).

An alternative model, which is arguably the most natural one, is a radiative contribution from the photosphere (Eichler & Levinson 2000; Mészáros & Rees 2000; Daigne & Mochkovitch 2002; Mészáros et al. 2002; Rees & Mészáros 2005; Ghirlanda et al. 2007; Beloborodov 2010; Ioka 2010; Mizuta, Nagataki & Aoi 2010). Indeed, a photospheric emission is a natural outcome of the ‘fireball’ model. At small radii, close to the inner engine, the optical depth is huge, \( \tau \sim 10^{15} \) (see e.g. Piran 2005, for a review). As a result, photons cannot escape, but are advected with the flow until they decouple at the photosphere. The huge value of the optical depth implies that regardless of the initial emitted spectrum, if photons are emitted deep enough in the flow, the spectrum emerging at the photosphere is blackbody (a Wien spectrum may be obtained if the number of photons is conserved). In fact, in the classical ‘fireball’ model, photons serve as mediators for energy conversion (the explosion energy is converted into kinetic energy of the relativistically expanding plasma jet by scattering between photons and leptons, which in turn convert the energy to the baryons via Coulomb collisions). Therefore, the appearance of thermal photons as part of the observed spectrum is not only expected, but is in fact required.

Apart from being an inherent part of the ‘fireball’ model, one of the key advantages of the idea that at least part of the observed spectrum originates from the photosphere is that it provides a relief for the efficiency problem of the internal shock model (Pe’er, Mészáros & Rees 2005; Lazzati, Morrisy & Begelman 2009; Ryde & Pe’er 2009). This is because if indeed part of the photons that we see have photospheric origin, then the total energy seen in photons is higher than the energy released by the internal dissipation: only the non-thermal part of the spectrum is required to originate from energy dissipation above the photosphere. Thus, the dissipated energy may be smaller than the remaining kinetic energy, even though the total energy seen in the photon component is comparable to the kinetic energy. A second great advantage of this idea is that the Rayleigh–Jeans part of the thermal spectrum (or modification of it) can naturally account for low-energy spectral slopes much steeper than those allowed by the synchrotron theory, and are hence consistent with observations.

The observed GRB spectra are therefore expected to be hybrid, i.e. containing both thermal and non-thermal parts. These two parts are connected in a non-trivial way. As the non-thermal part originates from energy dissipation (one or more) occurring above the photosphere, it is naturally delayed with respect to the thermal photons. A complete decomposition of the thermal and non-thermal parts of the spectrum is in fact nearly impossible. Thermal photons serve as seed photons to Compton scattering by the energetic electrons, thereby affecting the non-thermal part as well. Even if the dissipation occurs high above the photosphere, thermal photons significantly contribute to the cooling of the electrons (Pe’er et al. 2005), therefore affecting the spectrum emitted by these electrons. As was shown by Pe’er, Mészáros & Rees (2006), under these conditions a flat energy spectrum (\( \nu F(\nu) \propto \nu^0 \)), resulting from multiple Compton scattering of the thermal photons, is naturally obtained for a large range of parameters. A similar conclusion was drawn in the case of a magnetized outflow under the assumption of slow heating due to continuous reconnection at all radii (Giannios 2006). These predictions were shown to be in qualitatively very good agreement with Fermi results (Toma, Wu & Mészáros 2010).

The non-trivial connection between the thermal and non-thermal components, combined with the fact that both components vary in time, makes it difficult to identify the thermal component. A breakthrough in identifying this component was made by Ryde (2004, 2005), who looked at time resolved spectra, thereby allowing an identification of the temporal evolution of the temperature. Repeating a similar analysis on a fairly large sample of bursts, it was shown by Ryde & Pe’er (2009) that the thermal component not only could be identified, but that both the temperature and flux of this component show repetitive temporal behaviour at late times: \( F_{\nu,0} \propto \nu^{-1} \) and \( T^\theta \propto \nu^{-2/3} \). This behaviour was found to be in very good agreement with the theoretical predictions (Pe’er 2008), thereby providing an independent indication for the correct identification of the thermal emission component. It should be pointed out here that the thermal emission is expected to appear as grey-body emission, composed of multicolour blackbody spectra. This results from the fact that at a given time interval an observer sees photons emitted from different radii and different angles, hence undergoing different Doppler shifts. A full analysis shows that the resulting low-energy spectrum (below the observed spectral peak, typically seen at sub-MeV) is a power law (Pe’er & Ryde 2011).
Once the thermal component is identified, it is relatively easy to use its properties to deduce the dynamics of the outflow. As opposed to the non-thermal part, whose emission radius is uncertain (there may be multiple emission radii), the emission radius of the thermal photons is defined to be at the photosphere. Thus, by studying the properties (thermal flux and temperature) of the thermal component, under the assumption of constant outflow velocity, it is possible to deduce the photospheric radius and the Lorentz factor of the flow (Pe’er et al. 2007). We note that this method is independent, and is complementary to the opacity argument method commonly used to constrain the Lorentz factor of the flow at the emission radius of the high-energy photons (Svensson 1987; Krolik & Piran 1991; Woods & Loeb 1995; Lithwick & Sari 2001). It has two main advantages to the opacity argument method. First, it provides a direct measurement of the Lorentz factor, rather than a lower limit. Secondly, it is independent of measurements of the variability time, which is highly uncertain.

While in the past a clear identification of a thermal component was difficult, the situation dramatically changed with the broadband spectral coverage enabled with the launch of *Fermi*. Out of 14 GRBs detected by the Large Area Telescope (LAT) until 2010 January, three show clear evidence for a distinctive high-energy (>MeV and up to the GeV range) spectral component, and another eight show marginal evidence for such a component (Granot 2010). Thus, it is natural to deduce that the low-energy photons (below and at the sub-MeV peak of the flux) have a different origin than the higher energy (LAT) photons.

Out of the three bursts that show clear evidence for a distinct high-energy spectral component, GRB 090902B may be the easiest to analyse, for two reasons. First, the low-energy part of the spectrum (the spectral peak) is very narrow, and both the low-energy and high-energy spectral slopes are very steep. Thus, any attempts of fitting the spectrum using the standard synchrotron-SSC model are rejected. On the other hand, the spectrum is easily fitted with a (multicolour) blackbody spectrum, plus an additional power law (Ryde et al. 2010; Zhang et al. 2011). Thus, to date, this burst is unique by the fact that a thermal component is so clearly pronounced in its spectrum. Secondly, the high-energy power-law component is spectrally distinct from the low-energy one. Using the opacity argument, the detection of a 33.4–GeV photon associated with this burst (Abdo et al. 2009a) implies that the emission radius of this photon must be much greater than the photosphere (see below). These two facts make this burst ideal for demonstrating how separation of the spectrum into thermal and non-thermal components enables one to deduce the physical conditions of the outflow at the emission sites of both the thermal and non-thermal photons. Moreover, as we will show below, one can use measurements of the non-thermal part, to remove some of the uncertainties that exist in measurements of the thermal component alone.

This paper is organized as follows. In Section 2, we provide the general properties of the combined thermal and non-thermal model. We show how one can combine the hydrodynamic information derived by studying the properties of the non-thermal part of the spectrum, to provide a comprehensive picture of the physical parameters at both emission sites. In Section 3 we demonstrate the use of our method by fitting the prompt emission spectrum of GRB 090902B, and deducing its physical properties. We summarize and conclude in Section 4.

### 2 GENERAL PROPERTIES OF A THERMAL AND NON-THERMAL MODELS

#### 2.1 Temperature, luminosity and constraints on the value of $\epsilon_e$

We consider a fireball wind of total luminosity $L$, expanding from an initial radius $r_0$ (which, for the sake of argument, can be assumed to be a few times the last stable orbit around the central black hole or the sonic radius; in any case, it is a few times the Schwarzschild radius around a non-rotating black hole; see further discussion at the end of Section 3.1). The initial blackbody temperature at $r_0$ is $T_0 = (L/4\pi r_0^2 s/c^3)^{1/4}$, where $c$ is the speed of light and $a$ is the radiation constant. As the optically thick (adiabatic) wind expands, the baryon bulk Lorentz factor increases as $\Gamma \propto r$, and the comoving temperature drops as $T' \propto r^{-1}$ (see e.g. Mészáros 2006, for a comprehensive review).\(^1\) As long as the wind remains optically thick, the acceleration continues until the plasma reaches the saturation radius, $r_s = \eta r_0$, above which $\Gamma$ coasts to a value equal to the dimensionless entropy, $\eta \equiv L/Mc^2$. Here, $M$ is the mass outflow rate.

The photospheric radius is the radius above which the flow becomes optically thin to scattering by the baryon-related electrons. Depending on the values of the free model parameters ($\eta$, $L$ and $r_0$), this radius can be above or below the saturation radius (see Mészáros et al. 2002). For the parameter values characterizing GRBs (see below), the photospheric radius is above the saturation radius, and is given by (Abramowicz, Novikov & Paczyński 1991; Pe’er 2008)

$$r_{ph} = \frac{L\sigma_T}{8\pi m_p c^3} = 5.8 \times 10^{15} L_{42} \eta_3^{1/3} \text{cm.}$$

Here and below, $\sigma_T$ is the Thomson cross-section, $m_p$ is the proton mass, and we use the convention $Q = 10^{54} Q_4$ in cgs units. The high values of the luminosity and the Lorentz factor chosen for the demonstration in equation (1) are for ease of comparison with the *Fermi* results of GRB 090902B (see Section 3 below).

Above the saturation radius, adiabatic energy losses (in the absence of dissipation) cause the temperature to drop as $T = T_0 (r/r_s)^{2/3}$. The observed temperature of photons emitted at the photosphere is therefore

$$T^{ob} = T_0 \left( \frac{r_{ph}}{r_s} \right)^{-2/3} = 3.6 \times 10^5 (1 + z)^{-1} L_{42}^{5/12} \eta_3^{1/6} \Gamma_8^{-1/6} \text{eV.}$$

where $z$ is the redshift. Note the very strong dependence of the observed temperature on the asymptotic value of the Lorentz factor, $\eta$: for high $\eta$, high values of the temperature are expected.

The observed photospheric thermal luminosity drops above the saturation radius as $L^{ob} = (L/2)(r/r_s)^{2/3}$, the greater part of the energy being in a kinetic form, $L_k \sim L/2$ (Mészáros & Rees 2000; Rees & Mészáros 2005).\(^2\) The non-thermal part of the spectrum results from dissipation of the kinetic energy. Part of the dissipated energy goes into accelerating electrons that radiate the non-thermal spectrum. Denoting by $\epsilon_d$ the fraction of kinetic energy that is dissipated and by $\epsilon_e$ the fraction of dissipated energy that is converted

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1. From here on, quantities measured in the comoving frame are primed, while unprimed quantities are in the observer frame.

2. Note that in fact at $r > r_{ph}$ the kinetic luminosity is slightly higher due to energy conversion from the photons above the saturation radius. While the full treatment will be given below, in the content of equation (3) this has little effect, and is omitted for clarity.
to energetic electrons, one obtains an upper limit on the ratio of non-thermal to thermal luminosity in the spectrum:

$$\frac{L_{\text{NT}}^{\text{ob}}}{L_{\text{Th}}^{\text{ob}}} \leq \epsilon_d \epsilon (\frac{r_{\text{ph}}}{r_\gamma})^{2/3} = 0.33 L^{2/3} \eta^{8/3} r_{\text{ph}}^{2/3} \epsilon_d \epsilon_e e^{-1}. \quad (3)$$

where the non-equality results from the fact that the electrons do not necessarily radiate 100 per cent of their energy. Equation (3) implies an interesting result: the higher the Lorentz factor, the more pronounced its thermal luminosity is expected to be. This result may be very significant given recent Fermi-LAT data, which show evidence for high values of the Lorentz factors in several bursts. We further point out that as $\eta$ increases, the saturation radius $r_\gamma$ increases, while the photospheric radius decreases. Thus, at high enough value of $\eta = \eta^\ast \equiv (L \sigma T/4\pi m_e^2 c^4 r_\gamma)^{1/2}$, $r_{\text{ph}} = r_\gamma$, and the ratio $L_{\text{NT}}^{\text{ob}}/L_{\text{Th}}^{\text{ob}}$ saturates to $\epsilon_d \epsilon_e$ (see Mészáros et al. 2002; Pe’er et al. 2007, for further details). Higher value of $\eta > \eta^\ast$ does not change this result.

2.2 Initial expansion radius, photospheric radius and Lorentz factor

As was shown by Pe’er et al. (2007), the outflow parameters, in particular the Lorentz factor $\eta$, the initial expansion radius $r_\gamma$ and the photospheric radius $r_{\text{ph}}$, can be inferred directly from studying the thermal component alone (for bursts with known redshift). For completeness of the analysis, we briefly repeat here the main arguments given by Pe’er et al. (2007).

The observed temperature of thermal photons emitted from the photosphere is $T_{\text{ph}}^{\text{ob}} \simeq 1.48\eta T(r_{\text{ph}}/1 + z)$. Due to Lorentz aberration, the ratio $(F_{\text{ph}}^{\text{ob}}/\sigma T_{\text{ph}}^{\text{ob}})^{1/2}$ is proportional to the photospheric radius divided by the Lorentz factor:

$$R \equiv \frac{F_{\text{ph}}^{\text{ob}}}{\sigma T_{\text{ph}}^{\text{ob}}} = \frac{(1 + z)^2}{\eta} \frac{r_{\text{ph}}}{d_\gamma} \frac{L}{L_{\text{Th}}^{\text{ob}}}. \quad (4)$$

Here, $\sigma$ is Stefan’s constant, $\xi$ is a geometrical factor of order unity, $d_\gamma$ is the luminosity distance and $F_{\text{ph}}^{\text{ob}} = L_{\text{ph}}^{\text{obs}}/4\pi d_\gamma^2$ is the observed thermal flux. Using the parametric dependence of $r_{\text{ph}}$ from equation (1) in equation (4), one obtains the asymptotic Lorentz factor:

$$\eta = \left(\frac{\xi(1 + z)^2}{2m_e c^3} R\right)^{1/2} \left(\frac{L}{L_{\text{Th}}^{\text{ob}}}ight)^{1/4}. \quad (5)$$

Combining this result with the equation for the observed thermal flux, $L_{\text{Th}}^{\text{ob}} = (L/2)(r_{\text{ph}}/r_\gamma)^{2/3}$, one obtains the initial expansion radius:

$$r_\gamma = \frac{d_\gamma R}{\xi(1 + z)^2} \left(\frac{2L}{L_{\text{Th}}^{\text{ob}}}ight)^{3/2}, \quad (6)$$

and the photospheric radius:

$$r_{\text{ph}} = \left[\frac{d_\gamma R \sigma T_{\text{ph}}^{\text{ob}}}{\xi(1 + z)^2/2m_e c^3}\right]^{1/4} \left(\frac{L}{L_{\text{Th}}^{\text{ob}}}ight)^{1/4}. \quad (7)$$

The values of $\eta$, $r_\gamma$, and $r_{\text{ph}}$ are thus fully determined by the observed quantities of the thermal emission, up to the uncertainty in the luminosity ratio $L/L_{\text{Th}}^{\text{ob}} \geq 1$. Furthermore, equations (5) and (7) imply that the values of $\eta$ and $r_{\text{ph}}$ are not very sensitive to the uncertainty in this ratio.

Constraining the ratio of the total luminosity released in the explosion to the luminosity emitted as thermal photons, $L/L_{\text{Th}}^{\text{ob}}$, can most easily be done if an independent measurement of the kinetic energy exists. Such measurements are provided by studying the emission during the afterglow phase (Wijers & Galama 1999; Frail et al. 2001; Freedman & Waxman 2001; Panaitescu & Kumar 2001; Berger, Kulkarni & Frail 2003; Bloom, Frail & Kulkarni 2003; Nysewander et al. 2009), which provides good estimates of the kinetic energy remaining after the prompt emission phase. Fortunately, such measurements become ubiquitous and are available for GRB 090902B (Cenko et al. 2011); see further discussion in Section 3 below.

Even if afterglow measurements do not exist, the luminosity ratio $L/L_{\text{Th}}^{\text{ob}}$ can still be constrained indirectly, in three independent methods. First, by fitting the non-thermal part of the spectrum, one obtains a constraint on $\epsilon_d \epsilon_e (L/L_{\text{Th}}^{\text{ob}})$ (see equation 3). Since $\epsilon_d \epsilon_e < 1$, a lower limit on $L/L_{\text{Th}}^{\text{ob}}$ is obtained. Secondly, variability time measurements (if they exist) can constrain the initial expansion radius, since $\delta t \geq r_{\text{ph}} c$, which, in turn, provides an upper limit on the ratio $L/L_{\text{Th}}^{\text{ob}}$ via equation (6). And finally, a high value of $L_{\text{Th}}^{\text{ob}}$ implies that the luminosity ratio should not be high, in order to avoid an energy crisis.

2.3 Constraint on the emission radius of the non-thermal photons

Observations of high-energy ($\gtrsim 10\text{GeV}$) photons by Fermi, are commonly used in the literature to constrain the Lorentz factors of GRB outflows, using the opacity argument. We point out though that the constraints set in the literature are often based on an additional assumption, that is the emission radius of the high-energy photons, $r_{\gamma^*}$, is connected to the Lorentz factor via $r_{\gamma^*} = 2\gamma^* c \delta t$, where $\delta t$ is the variability time of the inner engine activity. This assumption, while true in the internal shock model scenario, has two main drawbacks. First, it assumes an a priori knowledge of the variability in the Lorentz factor, namely $\Delta \eta \approx \eta$ (this assumption translates into the numerical coefficient), and secondly, it relies on an assumed knowledge of the physical variability time, $\delta t$, which is difficult to be measured accurately.

A different approach was suggested by Zhang & Pe’er (2009): by releasing the requirement $r_{\gamma^*} = 2\gamma^* c \delta t$, it was shown that the opacity argument can be used to provide general constraints in the $r_{\gamma^*} - \eta$ space. While this method does not provide directly the value of $\eta$ (or of $r_{\gamma^*}$), its main advantage is that it is not sensitive to the uncertainties mentioned above. Here, we take the arguments presented by Zhang & Pe’er (2009) one step forward. We first use the analysis of the thermal component to estimate the Lorentz factor $\eta$. At the second step, we use the constraints found by the opacity argument in the $r_{\gamma^*} - \eta$ plane to deduce a lower limit on the emission radius of the non-thermal photons, $r_{\gamma^*}$.

The calculation is performed as follows. The cross-section for pair production of photon with energy $\epsilon_1$ is the highest for interactions with photons of energy $\epsilon_2 = (m_e c^2)^2/\epsilon_1$. Therefore, considering the Lorentz boosting, a photon observed with energy $\epsilon_{\text{obs}} = 10 \epsilon_{\text{obs max},10} \text{GeV}$ interacts mainly with photons having energies

$$\epsilon_{\text{obs}} \leq \frac{(m_e c^2)^2 \eta^2}{\epsilon_{\text{obs max}} (1+z)} \leq \frac{26}{(1+z)^2} \eta^2 \epsilon_{\text{obs max},10}^{-1} \text{MeV}. \quad (8)$$

This energy is about 2 orders of magnitude higher than the energy of the thermal photons (see equation 2). We therefore do not expect the thermal photons to play a significant role in constraining the
emission radius of the most energetic photons seen by the *Fermi*-LAT.

The observed spectrum above a few MeV is often modelled by a single power law, $dN/dE = f_0(E/E_0)^{-\gamma}$. For such a spectral fluence, the optical depth for pair production can be written as (Krolik & Pier 1991; Woods & Loeb 1995; Lithwick & Sari 2001; Zhang & Pe`er 2009)

$$\tau_{\gamma\gamma} = \frac{\langle \sigma \rangle d_4^2}{z_2^4 (1+z)^2} f_0 \Delta \epsilon_{\text{GeV}} \langle \epsilon_{\text{max}} \rangle^{1-\alpha} \left( \epsilon_{\text{min}} \right)^{1+\gamma} 1-2\epsilon \right)^2. \tag{9}$$

Here, $\langle \sigma \rangle$ is the cross-section averaged over all angles; for flat energy spectrum ($\alpha = 2$), one obtains $\langle \sigma \rangle \approx \sigma/8$ (Svensson 1987; Gupta & Zhang 2008). Further note that $\Delta \epsilon_{\text{GeV}}$ in equation (9) represents the observed time bin during which high-energy photons are seen, and thus does not necessarily correspond directly to the uncertain physical variability time.

For a flat energy spectrum ($\alpha = 2$) observed between $\epsilon_{\text{min}}$ and $\epsilon_{\text{max}}$, the observed fluorescence is related to the (non-thermal) luminosity via $f_0 = L_{\text{ob}}/4\pi d_4^2 \log(\epsilon_{\text{max}}/\epsilon_{\text{min}})$. The requirement that the optical depth to pair production of the highest energy photon seen is smaller than unity is translated into a lower limit on the emission radius,

$$\epsilon_{\text{min}} \geq \frac{\langle \sigma \rangle L_{\text{ob}}}{4\pi \log(\epsilon_{\text{max}}/\epsilon_{\text{min}}})^{1/2} \epsilon_{\text{min}}^{1/2} \eta_3^{-1} \text{ cm}. \tag{10}$$

where $\log(\epsilon_{\text{max}}/\epsilon_{\text{min}}) \approx 20$ was taken. For parameters characterizing GRBs seen by the LAT, the emission radius of the non-thermal photons is about 3–4 orders of magnitude larger than the photospheric radius (equation 1), indicating that the observed spectrum must be emitted from (at least) two separate regions.

### 2.4 The observed non-thermal spectrum: additional constraints on the free model parameters

The dissipation at $r_w$ regardless of its exact nature, produces a population of energetic electrons. The energetic electrons emit the non-thermal part of the spectrum, by radiating their energy. There are three main radiative mechanisms responsible for the non-thermal emission: synchrotron emission, SSC emission and Comptonization of the thermal photons.

In order to estimate the relative contributions of the different emission mechanisms to the observed spectrum, we proceed in the following way.

#### 2.4.1 Electron energy loss by Comptonization of the thermal photons

The thermal photons serve as seed photons for Compton scattering by the energetic electrons. Their existence therefore contributes to the high-energy part of the non-thermal spectrum.

The power emitted by Comptonization of the thermal photons (in the Thomson regime), relative to the power emitted as synchrotron radiation, is given by the ratio of the energy density of the thermal photon field to the energy density in the magnetic field. At the photosphere, the (comoving) energy density of the thermal photons is $u_\text{Th}(r_w) = aT(r_w)^4 = (L/\pi r_w^2 c^3)(r_w/r_d)^{-5/3}$. Since above the saturation radius the energy density of the photon field drops as $u_\text{Th} \propto r^{-2}$, at the dissipation radius $r_d$, it is equal to $u_{\text{Th}}(r_d) = (L/\pi r_d^2 c^3)(r_d/r_d)^{-5/3}$. The energy density in the magnetic field assumes a fraction $\epsilon_B$ of the comoving energy density, $L/8\pi r_d^2 c^3$. One therefore concludes that the ratio of the energy densities in the thermal photon and magnetic field is given by

$$\frac{u_{\text{Th}}}{u_B} = \frac{1}{\epsilon_B} \left( \frac{r_w}{r_d} \right)^{-2/3} = 3 \left( \frac{L}{r_d^2} \right) \frac{\eta_3^{3/2} \epsilon_{\text{ob}}^{2/3}}{r_d^{5/2}} \epsilon_B^{-1/2}. \tag{11}$$

It is thus clear that for parameters characterizing GRBs, the role played by Comptonization of thermal photons is, at the least, comparable to the role played by the synchrotron emission as a source of energy loss of the energetic electrons.

#### 2.4.2 Synchrotron spectrum

The dissipation process is expected to produce a power-law distribution of energetic electrons with power-law index $p \approx 2.0$, above a characteristic Lorentz factor $\gamma_{\text{min}} \approx (m_e c/\eta_5) \epsilon_{\text{ob}} \approx 184 \epsilon_{\text{ob}}^{-1}$. Comparison of the cooling time to the dynamical time implies that the entire electron population is in the fast cooling regime: the cooling time is equal to the dynamical time for electrons having Lorentz factor $\gamma_c = (3 m_e c^2 \eta_5 r_d) / (\eta_5 c L (1 + Y)) = 3.5 (1 + Y + Y^{-1})^{-1} L_{54}^{-1} \eta_3^{-1} \epsilon_{\text{ob}}^{-1}$. Here, $Y$ has its usual meaning as the ratio of SSC to synchrotron radiated power, $Y \equiv \eta_{\text{SSC}}/\eta_{\text{syn}}$.

The peak of the synchrotron emission, $\epsilon_{\text{ob}} = (3/2) \eta_5 (1 + z) q B r_s^2 / (m_e c) = 1.5 (1 + z)^{-1} L_{54}^{-1} \eta_3^{-1} \epsilon_{\text{ob}}^{-1}$ keV, is below the threshold energy of the *Fermi*-Gamma-ray Burst Monitor (GBM) detector. Here, $q$ is the electron charge and $B$ is the magnetic field at the dissipation radius. At the other end of the energy spectrum, comparison of the cooling time to the acceleration time provides an estimate of the maximum Lorentz factor of the accelerated electrons (assuming high efficiency in the acceleration process), $\gamma_{\text{max}} = [(\eta_5 c / \sigma_b) B (1 + Y + Y^{-1})]^{1/2}$. Synchrotron photons emitted by these electrons are expected at energies $\epsilon_{\text{ob}}^{\text{max}} = 240 (1 + z)^{-1} \eta_3^{-1} (1 + Y + Y^{-1})^{-1}$ GeV, above the threshold energy for pair production, and above the maximum photon energy seen so far by *Fermi*. These results imply that synchrotron emission is expected to contribute to the spectrum at the entire spectral range covered by *Fermi*.

The fast cooling of the electrons imply that (i) virtually all of the dissipated energy given to the electrons is radiated, and (ii) for power-law index $p \approx 2.0$, a flat energy spectrum ($\nu F_\nu \propto \nu^0$) is expected from synchrotron emission over the entire energy range covered by *Fermi*. Note though that a high-energy cut-off resulting from pair production can limit the maximum observed photon energy to values lower than $\epsilon_{\text{ob}}^{\text{max}}$ (see Section 2.3 above).

#### 2.4.3 Comptonization

There are two sources of Comptonized spectrum: Comptonization of the thermal photons and SSC. At low energies, below the thermal peak, Comptonization is not a significant source of photons. Hence, the spectrum below the thermal peak is dominated by synchrotron
emission. However, above the thermal peak, the three emission mechanisms — synchrotron, SSC and Comptonization of the thermal photons — contribute in parts to the spectrum. As shown in equation (11) and is further discussed below, the relative contributions of the different emission mechanisms are of the same order of magnitude (in other words, both \( T \) and \( Y \) are of the order unity). As a result, it is difficult to determine a single dominant emission mechanism at high energies. A consequence of this is that the observed spectral index at high energies (at the LAT band, above the thermal peak) cannot be directly related to a power-law index of the energetic electrons.

Comptonization of the thermal emission. Since \( (\gamma_y T)/m_e c^2 = 0.13 L_{54}^{-1/2} \eta_y^{-1/2} r_{0,8}^{-1/2} \epsilon_{-1} \approx 1 \), Comptonization of the thermal emission is in the Thomson limit. The Comptonized thermal photon spectrum has characteristic breaks similar to the SSC spectrum. Since the electrons are in the fast cooling regime, the Comptonized spectrum is expected to rise above \( \gamma_y^2 \Gamma^{\text{Th}} \gtrsim \Gamma^{\text{Th}} \) up to a peak at \( \epsilon_{\text{IC, obs}} = (4/3) \gamma_y^2 (2.87^{\text{Th}}) \approx 50 (1 + z)^{-1} L_{54}^{-1/2} \eta_y^{3/2} r_{0,8}^{-1/2} \epsilon_{-1} \) GeV, roughly as \( \approx \nu_F \propto \nu^{1/2} \) (Sari & Esin 2001). At the highest energies, photons annihilate by producing pairs, a phenomenon which can explain the lack of detection of the 50 GeV photons so far.

SSC. The ratio of SSC to synchrotron luminosity is given by the parameter \( Y \). For \( \epsilon_s \gg \epsilon_e \), the value of \( Y \) can be approximated as \( \approx (\epsilon_s/\epsilon_e)^{1/2} \) (Sari & Esin 2001).8 A significant high-energy non-thermal part, as is seen by the Fermi-LAT, implies (via equation 3) a high value of \( \epsilon_e \), close to equipartition. However, value of \( \epsilon_e > 10^{-2} \) as is inferred in many GRBs (see below), guarantees a value of \( Y \) of a few at most.

The SSC spectrum rises as \( \nu F_\nu \propto \nu^{1/2} \) below the peak of the SSC emission, which is expected at \( \epsilon_{\text{IC, obs}} = 2 \epsilon_m^{-2} \epsilon_s^{-1} \approx 100 (1 + z)^{-1} L_{54}^{-1/2} \eta_s^{-1/2} \epsilon_{-1}^{-1} r_{15,1}^{-1} \text{MeV} \). At higher energies, the spectral shape follows a similar power law as the synchrotron spectrum, i.e. a flat energy spectrum for \( p \approx 2 \).

The rise parts of both the thermal Comptonization and the SSC emission are independent of the power-law index of the accelerated electrons. For power-law index \( p > 2.0 \), a flat or slightly decaying synchrotron (energy) spectrum is expected at the entire Fermi energy range (see Section 2.4.2 above). The combined effects of the flat synchrotron emission with the rise of the Comptonized spectrum therefore results in a mild increase in the total observed spectrum at high energies, above the thermal peak (see Section 3 below). Demonstration of spectral decomposition into its basic physical ingredients is presented in Appendix A.

The rising of the SSC component below \( \sim 100 \text{MeV} \), combined with the fact that the synchrotron spectrum is expected to be flat (or slightly inverted) and that \( Y \) is not expected to be much larger than a few, implies that at low energies (below the thermal peak), Comptonization is not expected to play a significant role. The main emission mechanism below the thermal peak is therefore synchrotron emission. Since the Fermi-GBM detection range is above \( \epsilon_m \), the spectrum below the thermal peak is expected to be sensitive to the power-law index of the accelerated electrons, \( p \). Thus, measurement of the flux at low energies can provide an indication for both the values of \( \epsilon_g \) and \( p \).

Finally, we point out that the analysis carried in this section holds only as long as \( r_p \gg r_{\text{ph}} \). For \( r_p \gtrsim r_{\text{ph}} \), the effect of Comptonization is much more complicated due to the fact that the electrons cooling time is \( \propto r_p^{-1} \). Therefore, for small dissipation radius, the cooling time is much faster than the dynamical time for all electron energies, and \( \gamma_e \gtrsim 1 \). For such rapid cooling, additional physical phenomena that are not considered here become important. Direct Compton scattering of the energetic photons provide the main source of heating, resulting in a quasi-steady state distribution of mildly relativistic electrons (Pe’er et al. 2005). Close to the photosphere, multiple Compton scattering by electrons in this quasi-steady state produces a flat energy spectrum for a large region of parameter space (Pe’er et al. 2006).

3 Demonstration of the Analysis Method: GRB 090902B as a Concrete Example

The bright, long GRB 090902B, which is one of the brightest bursts observed by LAT to date (Abdo et al. 2009a), provides an excellent example for demonstrating our analysis method. This is because of two very pronounced properties of its prompt emission spectrum. First, time resolved spectral analysis reveals a significant power-law component in the LAT data (emission was observed up to 30 GeV), which is clearly distinct from the usual ‘Band’ function used by Abdo et al. (2009a) to fit the spectrum in the sub-MeV range (Ryde et al. 2010; Zhang et al. 2011). Moreover, the fact that the high-energy photons were delayed with respect to the photons at the sub-MeV peak indicates a different origin. Secondly, the ‘Band’ function used in fitting the sub-MeV peak is extremely steep on both sides (low-energy spectral slope \( n(\epsilon) \propto \epsilon^\alpha \), with \( \alpha > -0.5 \) to \(-0.3 \), and high-energy spectral slope \( \beta \approx -3 \) to \(-5 \), resulting in an unusually narrow peak. The steep spectral slopes seen in the sub-MeV range make it impossible to fit this spectrum with a model that contains only synchrotron and SSC.

On the other hand, both properties of the prompt spectrum fit perfectly into the framework suggested here. First, the sub-MeV peak is naturally modelled with the thermal emission component. While a single blackbody provides a sufficient fit, and is used to deduce the values of \( \eta, r_0 \) and \( r_{\text{ph}} \), better fits are obtained with multicolour blackbody, as expected from a theoretical point of view (Pe’er & Ryde 2011). Secondly, the non-thermal part (the high-energy power law which extends to lower energies) can be easily explained by a combination of synchrotron, SSC and Comptonization of the thermal photons. Moreover, by doing so, we deduce the physical properties in the emission sites of both the thermal and non-thermal components, hence we obtain a comprehensive physical picture of the properties of this burst.

In the original analysis, Abdo et al. (2009a) separated the observed prompt emission into several time bins. The most restrictive constraints on the emission radius of the non-thermal photons arise in time interval (c), 9.6–13.0 s from the trigger, since in this time interval an 11.16–GeV photon was observed. We therefore focus our analysis on this time interval.

3.1 Analysis of the thermal component: Lorentz factor, photospheric radius and dissipation radius

In order to deduce the value of the Lorentz factor by using the method presented in Section 2.2, one needs to fit the sub-MeV peak with a single blackbody spectrum, and study its properties (temperature and thermal flux). While this can be done for the data in time interval (c), since the properties of the outflow vary on a shorter time-scale (variability time of \( \lesssim 0.1 \) s was observed), smearing of the blackbody spectrum is expected. The shorter the time interval used for the fits, the higher the quality of the blackbody fits obtained (Ryde et al. 2010; Zhang et al. 2011). On the other hand,

8 Note that if \( Y = \tilde{Y} \), one obtains \( \gamma \approx (\epsilon_g/2\epsilon_e)^{1/2} \).
shorter time intervals result in lower quality of the high-energy data, which are sparse.

We thus use time interval as short as possible to fit the narrow peak with a single blackbody, in order to deduce the hydrodynamics of the flow, and in particular obtain the Lorentz factor. We then use the longer time interval [the full time interval (c)] to study the properties of the non-thermal part of the spectrum. For the long time interval, we fit the spectral peak with a multicolour blackbody, which, as was shown by Ryde et al. (2010), provides better fits to the peak. Clearly, by choosing to fit the longer time interval during which the values of the parameters vary, we lose an accuracy in the fits. However, we stress here that our purpose in this paper is not to provide the best statistical fits (in terms of $\chi^2$) to the data. Instead, our goal here is to prove that one can obtain an acceptable fits to the data in the entire Fermi energy band, in the sense that the curves plotted all fall within the $\pm 1\sigma$ error bars of the data points. By doing so, we show that we are able to provide an acceptable, complete physical interpretation to the observed data, although, necessarily, we are not able to capture many second-order effects.

For the shorter time interval we use the results presented by Ryde et al. (2010), who showed that the narrow time interval, 11.008–11.392 s from the trigger, during which the 11.16-GeV photon is seen, provides sufficient data to analyse the sub-MeV spectral peak. During this time interval, the sub-MeV peak can be fitted with a single blackbody, with observed temperature $T^{\text{ob}} = 168$ keV, and observed thermal flux $F^{\text{ob}} = 1.96 \times 10^{-5}$ erg cm$^{-2}$ s$^{-1}$. Equation (4) thus implies a ratio $R = 1.55 \times 10^{-19}$ cm. At redshift $z = 1.822$ (Cucchiara et al. 2009), the (isotropic-equivalent) thermal luminosity is $L^{\text{Th}}_{\text{ob}} = 4\pi d_L^2 F^{\text{ob}} = 4.6 \times 10^{53}$ erg s$^{-1}$. Using $\xi = 1.06$ in equations (5), (6) and (7) (Pe'er et al. 2007), one obtains $\eta = 764 (L/L^{\text{Th}}_{\text{ob}})^{1/4}$; $r_{\text{ph}} = 6.0 \times 10^{11} (L/L^{\text{Th}}_{\text{ob}})^{1/4}$ cm and $r_{\gamma} = 7.9 \times 10^{6} (L/2L^{\text{Th}}_{\text{ob}})^{1/2}$ cm.

Estimating the ratio $(L/L^{\text{Th}}_{\text{ob}})$ is done in the following way. Due to the rapid cooling of the electrons, nearly 100 per cent of the energy that is converted to the energetic electrons in the dissipation process is radiated in the form of non-thermal photons, i.e., $L^{\text{NT}}_{\text{ob}} \simeq \epsilon_e \epsilon_d L_{\text{k}}$. Here, $L_{\text{k}}$ is the energy available in kinetic form above the photosphere. By fitting the non-thermal part of the spectrum, one obtains $L^{\text{NT}}_{\text{ob}} \simeq 0.9 L^{\text{Th}}_{\text{ob}}$, or $\epsilon_e \epsilon_d \sim 0.9 (L^{\text{Th}}_{\text{ob}}/L_{\text{k}})$ (see Fig. 1). These fits are best done using a numerical simulation, since part of the dissipated energy is released outside the observed energy range of Fermi. After the main dissipation, the remaining kinetic luminosity, $L_{\text{AG}} = L_{\text{k}} (1 - \epsilon_d)$, is the available luminosity for the afterglow emission phase. As was found by Cenko et al. (2011), the energy released

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\caption{The dependence of the non-thermal flux on the fraction of energy given to the electrons, $\epsilon_e$. The data of GRB 090902B are from the NaI (0,1), BGO (0,1) and LAT (back and front) detectors at time interval (c), 9.6–13.0 s after the GBM trigger. The light yellow shaded area shows the $\pm 1\sigma$ fit to the data at this time interval, made by Abdo et al. (2009a). Note that this area is calculated by Abdo et al. (2009a) by fitting a ‘Band’ plus a single power-law spectra, hence the apparent discrepancy between the high-energy (LAT) data points and the shaded area. In our work, the parameter values chosen are dissipation radius $r_{\gamma} = 10^{17}$ cm, bulk motion Lorentz factor $\Gamma = 910$, power-law index of the accelerated electrons $p = 2.0$, GRB luminosity $L = 9.2 \times 10^{53}$ erg s$^{-1}$ and fraction of dissipated kinetic energy $\epsilon_d = 0.9$. Shown are the simulation results for $\epsilon_e = 0.5$ (blue, solid line) and $\epsilon_e = 0.17$ (dashed, green line). The $\sim$ MeV peak is fitted with a (multicolour) blackbody spectrum; the non-thermal flux is linear in the value of $\epsilon_e$.}
\end{figure}
during the afterglow phase is \( \sim 5 \) times less than the energy release during the prompt emission phase. However, since the luminosity in time interval (c) is about twice the average luminosity during the prompt phase, we estimate the luminosity (thermal + non-thermal) in time interval (c) to be about 10 times higher than the luminosity during the afterglow phase,

\[
\frac{L_{\text{AG}}}{L_{\text{Th}}^\text{ob} + L_{\text{NT}}^\text{ob}} = \frac{L_t(1 - \epsilon_d)}{L_{\text{Th}}^\text{ob} + \epsilon_d L_{\text{Th}}^\text{ob} \epsilon_d} \approx \frac{1}{10}
\]  

Using \( \epsilon_d \epsilon_d \approx 0.9(L_{\text{Th}}^\text{ob}/L_k) \), one obtains the relation:

\[
\epsilon_d = \frac{9}{1.9 \epsilon_d + 9}
\]  

(13)

This result immediately implies very high dissipation efficiency, since for any value of \( \epsilon_d \), \( \epsilon_d \geq 0.83 \).

At the photosphere, the thermal luminosity is equal to \( L_{\text{Th}}^\text{ob} = (L/2)(r_{\text{ph}}/r_c)^{-2/3} \), and the kinetic luminosity is equal to \( L_k = (L/2)(2 - (r_{\text{ph}}/r_c)^{-2/3}) \). At larger radii, the kinetic luminosity is assumed unchanged, up until the dissipation radius. Using the observed relation \( L_{\text{Th}}^\text{ob}/L_{\text{Th}}^\text{ob} = \epsilon_d L_k/L_{\text{Th}}^\text{ob} \approx 0.9 \) and the results obtained in equation (13), one obtains the luminosity ratio:

\[
\frac{L}{L_{\text{Th}}^\text{ob}} \approx \frac{(1.19 \epsilon_e + 0.9)}{\epsilon_e}
\]  

(14)

For equipartition value, \( \epsilon_e = 0.33 \), one obtains \( L/L_{\text{Th}}^\text{ob} = 3.9 \), which imply \( \eta = 1070 \) (and \( \epsilon_d = 0.94 \)). However, we consider this value of the Lorentz factor as an upper limit, due to two reasons: first, this result implies very high total GRB luminosity, \( L = 1.8 \times 10^{54} \text{erg s}^{-1} \); and secondly, using equation (6), it implies an initial expansion radius \( r_0 \approx 3 \times 10^5 \text{cm} \), which translates into very short variability time \( \delta t = r_0/c = 10 \text{ ms} \). Lower value of \( \epsilon_e \) results in higher Lorentz factor, but also leads to significantly higher total luminosity, significantly lower variability time and nearly 100 per cent dissipation efficiency, which we consider unlikely.

Using the extreme value \( \epsilon_e = 1 \), one obtains \( L/L_{\text{Th}}^\text{ob} = 2.09 \), \( \eta = 920 \), \( \epsilon_d = 0.83 \) and \( r_0 \approx 7.4 \times 10^5 \text{ cm} \), which translates into physical variability time \( \delta t \approx 25 \text{ ms} \). We thus conclude that the asymptotic value of the Lorentz factor is in the range \( 920 \leq \eta \leq 1070 \), and the value of \( \epsilon_e \) is at or slightly above equipartition. We further deduce that the photospheric radius is \( r_{\text{ph}} \approx 7.2-8.4 \times 10^{11} \text{ cm} \), and that the dissipation efficiency is \( \epsilon_d \approx 85-95 \) per cent. Using these values of \( \eta \) in equation (10) leads to the conclusion that the emission radius of the non-thermal photons is constrained, \( r_g \geq 3.5-4.1 \times 10^{13} \text{ cm} \).

The derived value of \( r_0 \approx 3 \times 7.5 \times 10^6 \text{ cm} \) implies \( r_0 \approx 33-80 r_{\text{ISCO},10} \) where \( r_{\text{ISCO},10} = 6 \text{ GM/c}^2 \) is the innermost stable circular orbit (ISCO) of a non-rotating 10-\( M_\odot \) black hole. As the exact mass of the progenitor of GRB 090902B is unknown, we can deduce that \( r_0 \) is of the order of few tens of the ISCO radius. According to the theory adopted here, \( r_0 \) marks the initial expansion radius. As the theory of jet acceleration is not fully developed yet, the values obtained may thus be used to constrain models of jet acceleration to this scale. Alternatively, we note that this value is very close to the value obtained in several numerical models (e.g. Aloy et al. 2000; Zhang, Woosley & MacFadyen 2003). While in these models the jets are assumed to be produced closer to the ISCO radius, the acceleration begins only at larger radii due to the fact that at smaller radii the jet is not well collimated.

### 3.2 Numerical calculations: further constraints on the free model parameters

As discussed in Section 2.4 above, following the dissipation process, simple analytical descriptions are insufficient to describe the spectrum, due to the roughly similar contributions from the different physical phenomena (synchrotron, SSC and thermal Comptonization), as well as the cut-off resulting from pair production. Therefore, in order to derive the spectral dependence on the different values of the free parameters, as well as confirm the analytical calculations presented above, we calculate numerically the photon and particle energy distribution for the different scenarios. In our calculations, we use the time-dependent numerical code presented in Pe’er & Waxman (2005). This code solved self-consistently the kinetic equations that determine the temporal evolution of electron and photons, describing cyclo-synchrotron emission, synchrotron self-absorption, direct and inverse Compton (IC) scattering, pair production and annihilation, and the evolution of high-energy cascade.

This code has two great advantages, which make it ideal in the study of the prompt spectra. First, it has a unique integrator, which enables solving the rate equations that govern the time evolution of the particles and photons energy distribution over the entire energy range. The code is able to calculate simultaneously processes happening over more than 15 orders of magnitude in time and energy ranges, thereby covering the entire spectral range, from radio up to the TeV band. The second advantage of the code is a full treatment of the various physical process, including the full cross-sections (e.g. Klein–Nishina effect is inherently taken into account, or that for mildly relativistic electrons, the full cyclo-synchrotron emission spectrum is calculated, etc.).

In modelling the spectrum in time interval (c), the sub-MeV peak is best described as a multicolour blackbody spectrum, with varying amplitudes (Ryde et al. 2010): \( F_{\nu}^\text{obs}(\nu; T_{\text{max}}) = \int_{\nu_{\text{min}}}^{\nu_{\text{max}}} dA(\nu)B_\nu(T) dA(\nu) \), where \( B_\nu(T) = (2\pi c^2)^{1/2}(\nu_{\text{obs}}^{3} - 1) \) is Planck function. The amplitude \( A(T) \) is temperature-dependent, \( A(T) = \frac{A(T_{\text{max}})(T_{\text{max}}/T)^{3-q}}{T_{\text{max}}^{1-q}} \), and is normalized such that the total flux is equal to the observed flux at the spectral peak, \( F_{\nu}^\text{obs} = 1.82 \times 10^{-5} \text{erg cm}^{-2} \text{s}^{-1} \). The normalization constant \( q \) and \( T_{\text{max}} \) are determined by fitting the spectrum, \( T_{\text{max}} = 328 \text{ keV} \) and \( q = 1.49 \). \( T_{\text{min}} \) cannot be determined, and its exact value is unimportant for the fits, as long as \( T_{\text{min}} \ll T_{\text{max}} \).

While the best value for \( q \) is found by fitting the data, we note that this multicolour description of the thermal part of the spectrum is inherent to emission from relativistically expanding plasma. At any given instance, an observer sees simultaneously photons emitted from a range of radii and angles (Pe’er 2008). The Doppler boosting of photons emitted at high angles to the line of sight is smaller than that of photons emitted on the line of sight. Therefore, a pure Planck function in the comoving frame is inevitably observed as multicolour blackbody.

In the framework developed here, the value of \( q \) is related to the spectral index via \( F_{\nu} \propto \nu^{q-1} \). This can be seen by noting that \( dA(\nu) \propto T^{q-5} \), and hence \( F_{\nu}^\text{obs}(\nu) = \int dA(\nu)B_\nu(T) \propto \nu^{q-1} \); the last equality is easily obtained by replacing the integrand from \( T \) to \( z = h\nu/T \). Thus, the value of \( q \) found by fitting the data implies spectral index \( F_{\nu} \propto \nu^{0.49} \). This index is significantly softer than the index expected for pure blackbody (\( F_{\nu} \propto \nu^1 \)), or from the scenario considered by Beloborodov (2010), of fixed comoving temperature, in which \( F_{\nu} \propto \nu^{1.4} \). On the other hand, this index is harder than the expected index at late times, \( F_{\nu} \propto \nu^{0.7} \) (Lundman, Pe’er & Ryde in preparation). The fitted spectral index is thus in between these values. We find this result encouraging, given that (i) here the fitted spectrum is integrated over a finite
Figure 2. The dependence of the non-thermal flux on the dissipation radius, $r_\gamma$. We show the numerical results for dissipation occurring at $r_\gamma = 10^{17}$ cm (solid, blue), $10^{16}$ cm (dashed, green), $10^{15.5}$ cm (dot-dashed, red) and $10^{15}$ cm (dotted, purple), superimposed on the Fermi data and the ±1σ ‘Band’ function fit to the data (light yellow shaded area). Values of $\epsilon_B = 0.33$ and $\epsilon_e = 0.5$ were chosen, apart from the fit for $r_\gamma = 10^{16}$ cm, where $\epsilon_e = 0.4$ is chosen. All the other parameters are the same as in Fig. 1. Below $10^{15.5}$ cm, pair production limits the maximum observed energy of photons to $\ll$ GeV, and is thus inconsistent with the LAT observation of 11.16-GeV photon at this time interval.

3.3 Numerical results

The numerical fits to the spectrum of GRB09092B at time interval (c), 9.6–13.0 s after the trigger, are presented in Figs 1–4. In Fig. 1, we demonstrate the linear dependence of the non-thermal flux on the value of $\epsilon_e$. For the fits shown in this figure, we chose parameter values that fulfil the requirements in the previous sections and are typical for GRBs. Thus, we chose large dissipation radius, $r_\gamma = 10^{17}$ cm, strong magnetic field, $\epsilon_B = 0.1$, and electron power-law index $p = 2.0$. The Lorentz factor was chosen to be $\Gamma = 910$, which implies $L/L_{\text{Th}} = 2.0$ or $L = 9.2 \times 10^{53}$ erg s$^{-1}$. These values imply, via equation (14), $\epsilon_e = 0.5$ and $\epsilon_d = 0.9$. The fit to the Fermi (GBM+LAT) data at time interval (c) appears as the solid (blue) line in Fig. 1. The fit shows the combined spectrum resulting from both the thermal peak (at $\sim$ MeV) and the nearly flat energy time interval, hence we average over spectra obtained at different times, and (ii) the theories are developed for the ‘pure’ case of constant outflow Lorentz factor. Hence, the results obtained are in good agreement with the theoretical expectations.9

In producing the spectrum, we assume that at radius $r_\gamma$, a fraction $\epsilon_d$ of the kinetic energy is being dissipated (by an unspecified dissipation process). The energetic electrons, which assume a power-law distribution with power-law index $p$, carry a fraction $\epsilon_e$ of the dissipated energy, and the magnetic field carries a fraction $\epsilon_B$ of this energy. The multicolour blackbody spectra serve as background spectra for all the various processes (mainly Compton scattering by the energetic electrons, but also other processes such as e.g. pair production). The code tracks the evolution of the spectrum, during the dynamical time. Here, we present the results at the end of time interval (c), i.e. we assume that the processes take place during an observed time of 3.4 s.

9 Additional broadening may result from sub-photospheric dissipation; see Ryde et al. (2011).

10 These can be compared to the spectra presented in fig. 3 in Abdo et al. (2009a), although note that in this work the spectrum is presented at a different time interval. However, the spectral shape and the main spectral features (such as the ratio of the flux at the peak to the flux at $\sim$10 keV) are similar.
The dependence of the non-thermal flux on the power-law index of the accelerated electrons. On top of the Fermi data, shown are the numerical results for $p = 2.0$ (solid, blue), $p = 2.2$ (dashed, green) and $p = 2.5$ (dot–dashed, red). Dissipation radius $r_{\gamma} = 10^{16}$ cm, $\epsilon_e = 0.5$, $\epsilon_B = 0.33$ and all other parameter values same as in Fig. 1 are chosen. The high-energy spectrum is nearly insensitive to the exact value of $p$ in the range considered, 2.0–2.5. However, the low-energy part (below the thermal peak) may provide indication for $2.0 \leq p \leq 2.2$.

We further note that the fit to the Wien part of the thermal component falls slightly below the $\pm 1\sigma$ error bars of the data (the shaded, yellow areas in the figures). This discrepancy can be easily understood as due to smearing of the data: the data presented in the figures is averaged over several seconds, during which the properties of the outflow (such as the Lorentz factor) slightly vary, while in the numerical fit we assume steady values of the physical parameters. Variation in the parameters values inevitably lead to smearing of the signal, which is translated to a high-energy decay which is somewhat shallower than the exponential cut-off of the thermal spectrum considered by the fits. Indeed, detailed analyses of time resolved spectra done by Ryde et al. (2010) and Zhang et al. (2011) show that as the time interval considered becomes shorter, the exponential decay above the thermal peak becomes more and more pronounced, hence the (multicolour) blackbody function used in fitting the peak becomes better the shorter the time bin is. None the less, as explained above, we chose here to fit the data in the entire time interval (c), since reducing the time interval results in poor quality of the high-energy part of the data (the non-thermal part). We therefore find it appropriate to use the fits presented, as they serve the main goal of this paper: to demonstrate that the physically motivated, hybrid (thermal+non-thermal) model provides acceptable fits, which, moreover, enable a good estimate of the physical conditions at both emission sites.

In Fig. 2, we consider different dissipation radii: $r_{\gamma} = 10^{17}$ cm (solid, blue), $10^{16}$ cm (dashed, green), $10^{15.5}$ cm (dot–dashed, red) and $10^{15}$ cm (dotted, purple). As the numerical code considers the full cross-section for pair production, the numerical results are more accurate than the analytical approximations presented in Section 2.3, and can be used to validate them. The results presented in Fig. 2 indeed confirm the main conclusion obtained analytically, that is that the observation of the 11.16-GeV photon necessitates the dissipation radius to be above $10^{15.5}$ cm.

At larger radii, the high-energy non-thermal part of the spectrum is not very sensitive to the exact dissipation radius. As shown in Fig. 2, for dissipation radii $r_{\gamma} \gtrsim 10^{15.5}$ cm it is possible to obtain numerical results which are within $\pm 1\sigma$ errors of the empirical ‘Band’ fit. In order to achieve this, high value of $\epsilon_B$ and a slight tuning of the value of $\epsilon_e$ is required. Thus, for $r_{\gamma} = 10^{16}$ cm, a value

**Figure 3.** The dependence of the non-thermal flux on the power-law index of the accelerated electrons. On top of the Fermi data, shown are the numerical results for $p = 2.0$ (solid, blue), $p = 2.2$ (dashed, green) and $p = 2.5$ (dot–dashed, red). Dissipation radius $r_{\gamma} = 10^{16}$ cm, $\epsilon_e = 0.5$, $\epsilon_B = 0.33$ and all other parameter values same as in Fig. 1 are chosen. The high-energy spectrum is nearly insensitive to the exact value of $p$ in the range considered, 2.0–2.5. However, the low-energy part (below the thermal peak) may provide indication for $2.0 \leq p \leq 2.2$. 

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Thermal and non-thermal emission in GRBs

477
Figure 4. The dependence of the non-thermal flux on the fraction of energy carried by the magnetic field, $\epsilon_B$. Shown are the numerical results for equipartition ($\epsilon_B = 0.33$; solid, blue), $\epsilon_B = 0.1$ (dashed, green) and $\epsilon_B = 0.01$ (dot–dashed, red). Dissipation radius $r_\gamma = 10^{16}$ cm, $\epsilon_e = 0.4$ and all other parameter values same as in Fig. 1 are chosen. While the effect of $\epsilon_B$ on the high-energy spectrum is minor, the low-energy flux (below the thermal peak) necessitates high value of $\epsilon_B$, close to equipartition.

of $\epsilon_e = 0.4$ was chosen, while for the other fits in this figure, $\epsilon_e = 0.5$ was found adequate.

The numerical results show that the dependence of the pair production cut-off at high energies on $r_\gamma$ is weaker than the analytical approximation presented in equation (10). The main reason for this is the contribution from the thermal photons, which is neglected in the derivation of equation (10). We can therefore conclude that further constraints on the dissipation radius cannot be obtained directly from the prompt spectrum, without additional assumptions. Finally, we point out that the small bump obtained in the scenario of $r_\gamma = 10^{15}$ cm at $\Gamma m_e c^2/(1 + z) \simeq 150$ MeV results from pair annihilation process, which is more pronounced at small dissipation radii due to the more rapid production of $e^\pm$ pairs. Such a small bump is difficult to be observed, and indeed has not been observed so far. Moreover, since the Lorentz factor of the flow likely varies during the time interval considered here, it is expected to be smeared.

In Figs 3 and 4 we examine the dependence of the spectra on the uncertain values of the power-law index of the energetic electrons and the fraction of dissipated energy carried by the magnetic field. In Fig. 3, we consider three values of the power-law index: $p = 2.0$ (solid, blue), $p = 2.2$ (dashed, green) and $p = 2.5$ (dot–dashed, red). As explained in Section 2.4.3 above, and is further demonstrated in Appendix A, the high-energy part of the spectrum (above the thermal peak) is governed by Compton scattering which is in the rising part of the spectrum (below $E_{IC}$), and therefore the spectrum is not sensitive to the exact power-law index of the accelerated electrons. We thus conclude that for GRB 090902B-type bursts, for which the thermal peak is pronounced, observations at high energies cannot be used to constrain the power-law index of the accelerated electrons. On the other hand, the low-energy part of the spectrum (below the thermal peak) is dominated by synchrotron emission, and is thus sensitive to the power-law index of the electrons. Unfortunately, most of the effect is expected below the threshold energy of the Fermi-GBM detector, and therefore only weak observational constraints exist. We can therefore conclude that a power-law index in the range $2.0 \leq p \leq 2.2$ is consistent with the data, and even the higher value of $p = 2.5$ can be consistent, however for such high power-law index a somewhat fine tuning of the other model parameters (in particular, very high value of $\epsilon_B$, close to equipartition) is required.

Examination of the spectral dependence on the value of $\epsilon_B$ is presented in Fig. 4. The three fits presented in this figure, equipartition ($\epsilon_B = 0.33$; solid, blue), $\epsilon_B = 0.1$ (dashed, green) and

11 Some other bursts such as e.g. GRB080916C show much less pronounced thermal peak, hence Comptonization may play a subdominant role for GRB080916c-type bursts.
\( \epsilon_B = 0.01 \) (dot–dashed, red), show that the value of \( \epsilon_B \) cannot be too far below equipartition: low value of \( \epsilon_B \) results in a too low flux at low energies (below the thermal peak), which is inconsistent with the observation. It also leads to a more pronounced Compton peak, which is seen at the high energies. We can thus conclude that value of \( \epsilon_B \) close to equipartition is needed to be consistent with the observations. This high value justifies the need for numerical analysis, since, although \( \epsilon_e \) is very high, both \( Y \) and \( \tilde{Y} \) are close to unity.

4 SUMMARY AND DISCUSSION

In this paper, we considered the effect of thermal emission on the observed GRB prompt emission spectrum. Being a natural outcome of the GRB fireball model, thermal emission is an inherent part of the prompt emission seen. As we showed in Section 2.1 (equation 3), it is expected to be more pronounced for bursts with higher Lorentz factor. As we show here, the inclusion of the thermal emission in the calculation of the prompt emission spectra enables one to obtain a complete, self-consistent physical model of the prompt emission spectrum seen over the entire Fermi energy range. This is in contrast to the ‘Band’ function fits, which do not carry any physical interpretation, and, in addition, require extra models to be able to fit the spectrum at high energies.

According to our model, the sub-MeV peak often seen is interpreted as being composed of multicolour blackbody emission from the photosphere. The non-thermal, high-energy part seen in several bursts by the LAT instrument is interpreted as resulting from combined emission of synchrotron, SSC and Comptonization of the thermal photons, following an episode of energy dissipation that occurs at large radius above the photosphere, \( r_p > r_{ph} \). Observations of high-energy photons can be used to constrain the dissipation radius (equation 10), which is clearly above the photospheric radius (equation 1). Moreover, as we showed in Section 2.4, the relative contributions of synchrotron emission, SSC and Comptonization of the thermal photons, denoted by the parameters \( Y \) and \( \tilde{Y} \), are expected to be of the same order of magnitude. Hence, the three emission mechanisms have roughly similar contributions to the high-energy (above the thermal peak) part of the spectrum. This fact makes it difficult to directly determine the power-law index of the accelerated electrons from measurements of the high-energy spectral slope.

The separation made here between thermal and non-thermal emission makes it possible to deduce the values of the free model parameters. First, by analysing the photospheric part one can deduce the value of the Lorentz factor, the initial expansion radius and the photospheric radius (see Section 2.2). Then, by analysing the non-thermal part one can constrain the dissipation radius, \( r_p \) (equation 10), and place constraints on the power-law index of the accelerated electrons, \( p \), and the strength of the magnetic field, \( \epsilon_B \). By combining the fluxes of the thermal and non-thermal parts, one can further constrain the combined fractions of dissipated kinetic energy (\( \epsilon_d \)) that is received by the energetic electrons (\( \epsilon_e \)). Further separation of the values of these variables is more tricky, but can be done with the help of afterglow observations (see Section 3.1, equations 12–14).

We demonstrated our analysis method on the bright, long GRB 090902B. This burst is ideal for our demonstration purposes, because of the clear separation between the thermal and non-thermal components seen, and the very pronounced thermal peak: both above and below the sub-MeV peak seen in this burst, the spectral slopes are too steep to enable fitting the spectrum with any combination of synchrotron and SSC emission models. However, as we showed in Section 3, Figs 1–4, excellent fits are obtained using the hybrid (thermal+non-thermal) model considered in this paper. While we stress again that we did not make any attempt to obtain the statistical best fits to the data, we are clearly able to obtain fits that are within the \( \pm 1 \sigma \) error bars of the data, over a very broad band – about 6 orders of magnitude spectral range. These fits are obtained using well-understood emission mechanisms. We can therefore conclude that, at least for this burst, the ‘Band’ function is well represented by a combination of physical emission processes. By doing so, we gain an insight into the physical conditions in the emitting regions.

According to our interpretation, the very pronounced \( \sim \)MeV peak represents the contribution of the thermal component, and the broad-band high- and low-energy spectra result from non-thermal processes. By making this separation, we were able to constrain the hydrodynamic properties of the outflow and the physical properties at the emission site: the initial expansion radius was found to be \( r_0 = 2.9–7.5 \times 10^8 \) cm, the photospheric radius \( r_{ph} \approx 7.2–8.4 \times 10^{11} \) cm and the Lorentz factor is in the range \( 290 \leq \eta \leq 1070 \). The main source of uncertainty in these values is the unknown kinetic luminosity.

Fit results from the interval during which an 11.16-GeV photon was observed constrain the dissipation radius to be \( r_p \geq 3.5–4.1 \times 10^{15} \) cm. Combined measurements of the thermal and non-thermal parts imply very high dissipation efficiency, \( \epsilon_d \approx 85–95 \) per cent, and very high fraction of dissipated energy carried by the energetic electrons: \( \epsilon_e \) is at or above equipartition value. Measurements of the low-energy part of the spectrum (below the thermal peak) imply high magnetic field \( \epsilon_B \) close to equipartition. The power-law index of the electrons is more difficult to constrain, as the high-energy spectral slope results from a combination of (flat or slowly decaying) synchrotron part, and rising Comptonized part. As such, the observed spectral slope does not directly correspond to the electron power-law index. From the numerical fits, we concluded that \( 2.0 \leq p \leq 2.2 \) are consistent with the data, and that even \( p = 2.5 \) is marginally consistent with the data.

The numerical results presented in Fig. 2 show high-energy cut-off due to pair production phenomenon. While this cut-off is consistent with the analytical prediction in equation (10), it also shows the limitation of the analytical approximations, which are commonly in use. Due to the inclusion of the thermal photons, the dependence of the cut-off energy on the emission radius is weaker than the simple approximation given in equation (10).

Additional constraints may be obtained by adding additional information, albeit with a higher level of uncertainty. For example, assuming the internal shocks scenario and considering that the 11.16-GeV photon was seen after a delay of 11 s from the GBM trigger implies, for constant Lorentz factor and \( \Delta t \approx \eta \), dissipation radius of \( r_p \sim \eta^2 \dot{c} t \approx 2 \times 10^{17} \) cm. We point out that this assumption is consistent with the observations, since the much more rapid variability time (\( \leq 0.1 \) s) can be attributed to emission from the photosphere, rather than the high-energy non-thermal part.

While GRB 090902B shows very pronounced peak and clearly separated high-energy component, which make it ideal for demonstrating our analysis method, this is not the case in many GRBs (see e.g. recent analysis by Zhang et al. 2011). In fact, in many GRBs, the sub-MeV peak often seen is not as pronounced as in GRB 090902B: the low- and high-energy spectral slopes are not as steep, and so in many cases pure single blackbody spectrum is too narrow to fit the sub-MeV peak. In some cases, e.g. GRB 100724B (Guiriec et al. 2011), a weak thermal component, which is not associated with the
main peak of the spectrum, can be identified. In many other cases, there is no clear evidence for the existence of thermal component, as blackbody spectrum is not clearly identified (Nava et al. 2011).\(^{12}\) Moreover, often the high-energy spectral slope (above the thermal peak) decays with a power-law index much shallower than seen in GRB 090902B. Thus, to date, the spectrum of GRB 090902B is unique by having such a pronounced thermal peak.

Clearly, a lack of very pronounced thermal peak in most bursts is a major drawback to the ideas raised here. There are several ways that can explain these observations. One possibility is suppression of the photospheric component which is expected in Poynting-flux-dominated flow (Zhang & Pe’er 2009; Zhang & Yan 2011). In this scenario, the magnetic field, rather than the photon field, serves as an energy reservoir. As the magnetic energy is gradually dissipated, most of the emission occurs at large radii above the photosphere, leaving only a weak photospheric signal. Within the fireball model itself, the irradiation of the thermal peak strongly depends on the value of the Lorentz factor \(\Gamma\), which is a free parameter of the model [see equations (2) and (3)]. While there is a clear indication for \(\Gamma \approx 10^2\) in GRB 090902B and several other bursts, in many bursts the value of \(\Gamma\) is lower, \(\sim 10^{2.5}\) (see e.g. Racusin et al. 2011). Thus, for these bursts, the thermal emission, while expected to exist, is not as pronounced as in GRB 090902B.

Alternatively, these observations can be explained in a framework similar to the one used here, namely that the spectrum may be dominated by Comptonization of the thermal photons, following energy dissipation that occurs close to the photosphere. As shown by Pe’er et al. (2005, 2006), in this scenario, multiple Compton scattering by electrons at a quasi-steady state produces a flat energy spectra for a large parameter space region. Since the dissipation radius \((r_x)\) can in principle take any value above the saturation radius, different dissipation radii can lead to very different observed spectra (see Pe’er et al. 2006); in particular, high-energy power-law tail is obtained (Lazzati & Begelman 2010). Thus, sub-photospheric dissipation can reproduce the Band function (Ryde et al. 2011). In fact, it could very well be that the uniqueness of GRB 090902B originate from a very large dissipation radius, \(r_x \gg r_{ph}\). Only under this condition, one is able to make such a clear separation between the thermal and non-thermal parts of the spectrum, which are otherwise coupled.

This possibility can also explain the lack of GeV emission in many bursts. If \(r_x\) is not much larger than \(r_{ph}\), then, by definition, the optical depth to scattering is high. Since the cross-section to pair production is similar to \(\sigma_T\), the optical depth to pair production is high too (see equation 10). As a result, GeV emission is attenuated. Alternatively, attenuation of GeV emission is expected if the outflow is highly magnetized: in this scenario, synchrotron emission dominates over IC scattering, leading to attenuation of the high-energy emission.

In addition to the complex relations between the thermal and non-thermal parts of the spectra, there are two effects which are often being neglected. First, the fits are often made to time-integrated spectra. As the properties of the outflow, in particular the Lorentz factor, vary on a very short duration [variation can be expected on time-scale of the order of \(r_0/c\), i.e. \(O(10\text{ms})\)], the blackbody spectrum is often smeared. Even more than that, as was shown in Pe’er (2008) and Pe’er & Ryde (2011), at any given instance, an observer sees simultaneously thermal photons emitted from different radii and different angles to the line of sight, hence having different Doppler shifts. As a result, the expected photospheric emission is not a pure blackbody, but a combination of blackbody spectra with different amplitudes. Thus, in fact, one expects to see a multicolour blackbody, as is indeed seen. The full theory of multicolour blackbody emission from relativistically expanding plasmas recently appeared in Pe’er & Ryde (2011).

The results of the fits to GRB 090902B imply a very high value of \(\epsilon\), close to or even above equipartition, and an even higher value of \(\epsilon_g\), 85–95 per cent. These values are much higher than the values predicted by the internal shock model: the typical efficiency in energy dissipation by internal shocks is no more than a few per cent. The results obtain here (note that the exact nature of the dissipation process is not specified) thus raise another issue as to the validity of the internal shock model scenario. This is added to GRB 080916C, in which detailed analysis by Zhang & Pe’er (2009) concluded that an additional source of energy must exist between the photospheric radius and the dissipation radius. The most plausible source of energy considered is magnetic, i.e. a Poynting-dominated outflow (see Zhang & Yan 2011). We can conclude that the high efficiency required may hint towards Poynting-dominated outflow in the case of GRB 090902B as well, and may even be a general requirement for all the bursts with pronounced high-energy emission observed by LAT. On the other hand, even if the outflow in GRB 090902B is Poynting-flux-dominated, we do not expect a too high ratio to kinetic luminosity, \(\sigma\), since \(\sigma \gg 1\) results in suppression of the photospheric emission, which is not observed in this burst. A more generalized treatment of photospheric models with arbitrary magnetization is outside the scope of this manuscript, and is left for future work.

The separation made in this work into two emission zones, namely thermal emission originating from the photosphere and non-thermal emission originating from energy dissipation at larger radii, provides a natural explanation to the delay of the high-energy photons, often seen in Fermi-LAT bursts (Abdo et al. 2009a; Ghisellini et al. 2010; Ryde et al. 2010; Toma et al. 2010). In our model, the non-thermal photons originate from dissipation above the photosphere, hence they are naturally seen at a delay with respect to the photospheric (thermal) photons, which are always the first to be observed. A pronounced thermal component at early times also provides a natural explanation to the harder slope at low energy observed during the first 1–2 s in many long bursts (Ghirlanda et al. 2009). In this general framework, the origin of the high-energy, non-thermal photons can also be hadronic, as recently suggested (Asano, Inoue & Mészáros 2010; Razzaque, Dermer & Finke 2010). None the less, we showed here that a leptonic origin is consistent with the data. Our model has the advantage that it does not require a large amount of energy in the hadronic component. Moreover, we did not specify the origin of the dissipation that led to the emission of the high-energy photons. Recently, there were several suggestions of external shock origin of these photons (Kumar & Barniol-Duran 2009, 2010; Ghisellini et al. 2010). However, by fitting the broadband data we showed that excellent fits can be obtained if both the characteristic synchrotron breaks are below the Fermi energy band, i.e. less than \(\sim 10\text{keV}\). Thus, our model is consistent with internal dissipation origin of the non-thermal spectrum (see further discussion in Piran & Nakar 2010; Zhang et al. 2011).

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\(^{12}\) Although the steepness of the low-energy spectral slope can be viewed as an indirect evidence.
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APPENDIX A: DECOMPOSITION OF THE SPECTRUM INTO ITS BASIC PHYSICAL INGREDIENTS

The observed spectrum results from synchrotron emission, SSC and thermalization of the thermal photons, with a high-energy cut-off resulting from pair production. As discussed in Section 3.2, the three emission processes are expected to have similar contributions to the observed spectrum of GRB 090902B. Therefore, a ‘clean’ decomposition into the spectral ingredients, as is presented in e.g. Sari & Esin (2001), does not exist in practice.

The high-energy spectral slope (above the thermal peak) does not have a direct correspondence to the power-law index of the

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accelerated electrons. Therefore, a decomposition of the spectrum into the basic radiative processes can be useful in understanding the physical processes that shape the spectrum. Fortunately, such a decomposition can (up to some level) be done numerically.

The numerical results are presented in Fig. A1, for dissipation radius $r_\gamma = 10^{16}$ cm and power-law index $p = 2.2$. The dot-dashed (red) curve represents the spectrum that would have been obtained if only synchrotron emission was considered. As the cooling frequency is low, the spectral index in the entire Fermi range is $\nu F_\nu \propto \nu^{-p/2} \propto \nu^{-0.1}$. When SSC is added (dashed, green curve) the spectral slope at low energies (below $\approx 10^4$ eV) is not affected, while the spectrum at high energies becomes nearly flat. This results from a combination of a decreasing synchrotron flux and an increasing SSC flux. Note that in this scenario the flux at low energies is weaker than the flux obtained for the pure synchrotron scenario, due to the fact that part of the electron energy is converted to SSC. Finally, the solid (blue) curve shows the combined effects when all the physical ingredients are added. Inclusion of the thermal photons does not affect the low-energy spectral slope (below the thermal peak). However, Comptonization of the thermal photons contribute to the more pronounced spectrum at high energies. In addition, the inclusion of thermal photons leads to a sharper high-energy cut-off, resulting from pair production.