Electron-positron plasma in GRBs and in cosmology

R. Ruffini\(^{(1)}\)(\(^{(2)}\)) and G. V. Vereshchagin\(^{(1)}\)(\(^{(2)}\))

\(^{(1)}\) Dipartimento di Fisica, Sapienza Università di Roma, ICRA - International Center for Relativistic Astrophysics, P.le Aldo Moro 5, Rome, 00185, Italy
\(^{(2)}\) ICRANET - P.zza della Repubblica 10, Pescara, 65122, Italy

Summary. — Electron-positron plasma is believed to play an important role both in the early Universe and in sources of Gamma-Ray Bursts (GRBs). We focus on analogy and difference between physical conditions of electron-positron plasma in the early Universe and in sources of GRBs. We discuss a) dynamical differences, namely thermal acceleration of the outflow in GRB sources vs cosmological deceleration; b) nuclear composition differences as synthesis of light elements in the early Universe and possible destruction of heavy elements in GRB plasma; c) different physical conditions during last scattering of photons by electrons. Only during the acceleration phase of the optically thick electron-positron plasma comoving observer may find it similar to the early Universe. This similarity breaks down during the coasting phase. Reprocessing of nuclear abundances may likely take place in GRB sources. Heavy nuclear elements are then destroyed, resulting mainly in protons with small admixture of helium. Unlike the primordial plasma which recombines to form neutral hydrogen, and emits the Cosmic Microwave Background Radiation, GRB plasma does not cool down enough to recombine.

PACS 52.27.Ep — Electron-positron plasmas.
PACS 52.27.Ny — Relativistic plasmas.
PACS 98.80.-k — Cosmology.

1. – Introduction

Electron-positron plasmas are discussed in connection with astrophysical phenomena such as Galactic Center, microquasars, Gamma-Ray Bursts (GRBs), as well as laboratory experiments with high power lasers, for details see [1]. According to the standard cosmological model, such plasma existed also in the early Universe. It is naturally characterized by the energy scale given by the electron rest mass energy, 511 keV. It is interesting that at the epoch when Universe had this temperature, several important phenomena took place almost contemporarily: electron-positron pair annihilation, the Big Bang Nucleosynthesis (BBN) and neutrino decoupling.

Electron-positron plasma also is thought to play an essential role in GRB sources, where simple estimates for the initial temperature give values in MeV region. Such
plasma is energy dominated and optically thick due to both Compton scattering and electron-positron pair creation, and relaxes to thermal equilibrium on a time scale less than \(10^{-11}\) sec, see [2]. The latter condition results in self-accelerated expansion of the plasma until it becomes either transparent or matter dominated.

In the literature there have been several qualitative arguments mentioning possible similarities between electron-positron plasmas in the early Universe and in GRB sources. However, until now there is no dedicated study which draws analogies and differences between these two cases. This paper aims in confronting dynamics and physical conditions in both cases.

2. – General equations

The framework which describes electron-positron plasma both in cosmology and in GRB sources is General Relativity. Both dynamics of expansion of the Universe, and the process of energy release in the source of GRB should be considered within that framework. Hydrodynamic expansion of GRB sources may, however, be studied within much simpler formalism of Special Relativity.

We start with Einstein equations

\[
R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = \frac{8\pi G}{c^4}T_{\mu\nu},
\]

where \(R_{\mu\nu}, g_{\mu\nu}\) and \(T_{\mu\nu}\) are respectively Ricci, metric and energy-momentum tensors, \(G\) is Newton’s constant, \(c\) is the speed of light, and the energy-momentum conservation, following from (1)

\[
(T_{\mu \nu})_{;\nu} = \frac{1}{\sqrt{-g}} \frac{\partial (\sqrt{-g} T_{\mu \nu})}{\partial x^\nu} - \Gamma^\lambda_{\nu \mu} T_{\lambda \nu} = 0,
\]

where \(\Gamma^\mu_{\nu \lambda}\) are Cristoffel symbols and \(g\) is determinant of the metric tensor. We assume for the energy-momentum tensor

\[
T^{\mu \nu} = p g^{\mu \nu} + \omega U^\mu U^\nu,
\]

where \(U^\mu\), is four-velocity, \(\omega = \rho + p\) is proper enthalpy, \(p\) is proper pressure and \(\rho\) is proper energy density.

When plasma is optically thick, radiation is trapped in it and entropy conservation applies. It may be obtained multiplying (2) by four-velocity

\[
-U^\mu (T_{\mu \nu})_{;\nu} = U^\mu \rho_{;\mu} + \omega U^\mu_{;\mu} = 0.
\]

Using the second law of thermodynamics

\[
d \left( \frac{\omega}{n} \right) = T d \left( \frac{\sigma}{n} \right) + \frac{1}{n} dp,
\]

where \(\sigma = \omega/T\) is proper entropy density, \(T\) is temperature, one may rewrite (4) as

\[
(\sigma U^\mu)_{;\mu} = U^\mu \sigma_{;\mu} + \sigma U^\mu_{;\mu} = 0.
\]
Baryon number conservation equation has exactly the same form

\[(nU^\mu)_{;\mu} = U^\mu n_{;\mu} + nU^\mu_{;\mu} = 0.\]

Now recalling that \[U^\mu \partial / \partial x^\mu = d / d\tau\] and \[U^\mu_{;\mu} = d \ln V / d\tau,\] where \(V\) is comoving volume, \(\tau\) is the proper time, from (4) and (7) we get

\[d\rho + \omega d\ln V = 0, \quad d\ln n + d\ln V = 0,\]

Finally, introducing the thermal index \(\gamma = 1 + \frac{p}{\rho}\) restricted by the inequality \(1 \leq \gamma \leq 4/3\) we obtain from (4) the following scaling laws

\[\rho V^\gamma = \text{const}, \quad nV = \text{const}.\]

Both these conservations laws are valid for the early Universe and GRB plasmas.

One can obtain the corresponding scaling laws for comoving temperature by splitting the total energy density into nonrelativistic (with \(\gamma = 1\)) and ultrarelativistic (with \(\gamma = 4/3\)) parts with \(\rho \rightarrow nmc^2 + \varepsilon\), where \(m\) is the mass of particles\(^{1}\), \(\varepsilon\) is proper internal energy density. The entropy of the ultrarelativistic component is then \(\sigma = \frac{4}{3} T\), and (6) gives

\[\frac{\varepsilon V}{T} = \text{const}.\]

For \(\varepsilon \gg nmc^2\), which is the energy dominance condition, internal energy plays dynamical role by influencing the laws of expansion. For \(\varepsilon \ll nmc^2\), which is the matter dominance condition, internal energy does not play any dynamical role, but determines the scaling law of the temperature. In order to understand the dynamics of thermodynamic quantities in both early Universe and in GRBs, one should write down the corresponding equations of motion.

21. Early Universe. – For the description of the early Universe we take the Robertson-Walker metric with the interval

\[ds^2 = -c^2 dt^2 + a^2(t) \left[ \frac{dr^2}{1 - kr^2} + r^2 d\theta^2 + r^2 \sin^2 \theta d\varphi^2 \right],\]

where \(a(t)\) is the scale factor and \(k = 0, \pm 1\) stands for the spatial curvature. In homogeneous and isotropic space described by (11), Einstein equations (1) are reduced to

\(^{1}\) Nonrelativistic component is represented by baryons. For simplicity we assume only one sort of baryons, say protons, having mass \(m\). Ultrarelativistic component is represented by photons and electron-positron pairs.
Friedmann equations together with the continuity equation

\begin{equation}
\left(\frac{da}{dt}\right)^2 + c^2 k = \frac{8\pi G}{3c^2} \rho a^2, \tag{12}
\end{equation}

\begin{equation}
2a \frac{d^2a}{dt^2} + \left(\frac{da}{dt}\right)^2 + c^2 k = -\frac{8\pi G}{c^2} \rho a^2, \tag{13}
\end{equation}

\begin{equation}
\frac{d\rho}{dt} + \frac{3}{a} \frac{da}{dt} (\rho + p) = 0, \tag{14}
\end{equation}

where \(a\) is the scale factor. Notice, that only two equations in the system above are independent. The continuity equation (14) follows from the Einstein equations (12) and (13) as the energy conservation. In fact, (14) may be also obtained from the entropy conservation (4). The comoving volume in Friedmann’s Universe scales with \(a\) as \(V = a^3\), so (14) and the first equality in (9) are equivalent.

On the radiation dominated stage of the Universe expansion one has

\begin{equation}
\rho \propto V^{-4/3} \propto a^{-4}, \quad n \propto V^{-1} \propto a^{-3}, \tag{15}
\end{equation}

while on the matter dominated stage

\begin{equation}
\rho \propto n \propto V^{-1} \propto a^{-3}. \tag{16}
\end{equation}

Entropy conservation (10) leads to the unique temperature dependence on the scale factor

\begin{equation}
T \propto V^{-1/3} \propto a^{-1}. \tag{17}
\end{equation}

The corresponding time dependence of thermodynamic quantities may be obtained from solutions of Friedmann equation (12) and continuity equation (14), see e.g. [3].

2’2. GRBs. – Different situation takes place for the sources of GRBs. Assuming spherical symmetry for the case of GRB the interval (2) is

\begin{equation}
ds^2 = -c^2 dt^2 + dr^2 + r^2 d\theta^2 + r^2 \sin^2 \theta d\varphi^2. \tag{18}
\end{equation}

Optically thick to Compton scattering and pair production electron-positron plasma in GRB sources is radiation dominated. Its equations of motion follow from the energy-momentum conservation law (2) and baryon number conservation law (7). Initially plasma expands with acceleration driven by the radiative pressure.

In spherically symmetric case the number conservation equation (7) is

\begin{equation}
\frac{\partial (n\Gamma)}{\partial t} + \frac{1}{r^2} \frac{\partial}{\partial r} \left( r^2 n \sqrt{\Gamma^2 - 1} \right) = 0, \tag{19}
\end{equation}

\(^{(2)}\) General Relativity effects may be included by taking Schwarzschild or Kerr-Newman metric. However, we are interested in optically thick plasma which expands with acceleration and propagates far from its source, where the spatial curvature effects may be neglected. For this reason we simplify the treatment and adopt a spatially flat metric.
Integrating this equation over the volume from certain $r_i(t)$ to $r_e(t)$ which we assume to be comoving with the fluid $\frac{dr_i(t)}{dt} = \beta(r_i, t)$, $\frac{dr_e(t)}{dt} = \beta(r_e, t)$, and ignoring a factor $4\pi$ we have

\begin{equation}
\frac{\partial}{\partial t} \int_{r_i}^{r_e} \frac{n_{\Gamma} r^2}{r \, \Gamma} dr + \frac{\partial}{\partial r} \left( r^2 n \sqrt{\Gamma^2 - 1} \right) \, dr = 0 \tag{20}
\end{equation}

\begin{equation}
\frac{\partial}{\partial t} \int_{r_i}^{r_e} (n\Gamma) r^2 dr - \frac{dr}{dt} n(r_e, t) \Gamma(r_e, t) r_e^2 + \frac{dr}{dt} n(r_i, t) \Gamma(r_i, t) r_i^2 +
\end{equation}

\begin{equation}
+ r_e^2 n(r_e, t) \sqrt{\Gamma^2(r_e, t) - 1} - r_i^2 n(r_i, t) \sqrt{\Gamma^2(r_i, t) - 1} = 0
\end{equation}

Since we deal with arbitrary comoving boundaries, this means that the total number of particles integrated over all differential shells is conserved

\begin{equation}
N = 4\pi \int_0^{R(t)} n\Gamma r^2 dr = \text{const}, \tag{21}
\end{equation}

where $R(t)$ is the external radius of the shell.

Following [4] one can transform (19) from the variables $(t, r)$ to the new variables $(s = t - r, r)$ and then show that

\begin{equation}
\frac{1}{r^2} \frac{\partial}{\partial r} \left( r^2 n \sqrt{\Gamma^2 - 1} \right) = - \frac{\partial}{\partial s} \left( \frac{n}{\Gamma + \sqrt{\Gamma^2 - 1}} \right). \tag{22}
\end{equation}

For ultrarelativistic expansion velocity $\Gamma \gg 1$, the RHS in (22) tends to zero, and then the number of particles in each differential shell between the boundaries $r_i(t)$ and $r_e(t)$ is also conserved with a good approximation, i.e.

\begin{equation}
dN = 4\pi n\Gamma r^2 dr \approx \text{const}. \tag{23}
\end{equation}

Relations (21) and (23) then imply

\begin{equation}
4\pi \int_{r_i}^{r_e} (n\Gamma r^2) dr = 4\pi \left[ n(r, t) \Gamma(r, t) r^2 \right] \int_{r_i}^{r_e} dr = 4\pi \left( n\Gamma r^2 \right) \Delta \approx \text{const}, \tag{24}
\end{equation}

where the first argument of functions $n(r, t)$ and $\Gamma(r, t)$ is restricted to the interval $r_i < r < r_e$, and consequently $\Delta \equiv r_e - r_i \approx \text{const}$. Taking into account that $r_i(t)$ and $r_e(t)$
are arbitrary, this means that ultrarelativistically expanding shell preserves its width measured in the laboratory reference frame. This fact has been used in [5] and referred there as the constant thickness approximation.

The volume element measured in the laboratory reference frame is \( dV = 4\pi r^2 dr \), while the volume element measured in the reference frame comoving with the shell is \( dV = 4\pi \Gamma r^2 dr \). Comoving volume of the expanding ultrarelativistic shell with \( \Gamma \simeq \text{const} \) will be

\[
V = 4\pi \Gamma \int_{r-\Delta}^{r} r^2 dr \simeq 4\pi \Gamma r^2 \Delta.
\]

Then we rewrite the conservation equations (9) as

\[
\rho \nabla \Gamma r^2 = \text{const}, \quad n \Gamma r^2 = \text{const},
\]

Unlike the early Universe, where both energy and entropy conservations reduce to (14), in the case of GRBs the energy conservation is a separate equation coming from the zeroth component of (2) as

\[
(T_0 \nu)_{,\nu} = \omega U_0 U^\nu \nabla U^\nu (\omega U_0)_{,\nu} = 0.
\]

which independently of \( \gamma \) gives

\[
\rho \Gamma^2 r^2 = \text{const}.
\]

From (26) and (28) we then find

\[
\Gamma \propto r^{\frac{2(\gamma-1)}{\gamma-2}}, \quad n \propto r^{-\frac{2}{\gamma-2}}, \quad \rho \propto r^{-\frac{2\gamma}{\gamma-2}}.
\]

For the ultrarelativistic equation of state with \( \gamma = 4/3 \) we immediately obtain

\[
\Gamma \propto r, \quad n \propto r^{-3}, \quad \rho \propto r^{-4}.
\]

Taking into account that the relation between the comoving and the physical coordinates in cosmology is given by the scale factor \( a \), it follows from (30) that both energy density and baryonic number density behave as in the radiation dominated Universe, see (15). This analogy between the GRB source and the Friedmann Universe is noticed by [6], [4].

In the presence of baryons as the pressure decreases, plasma becomes matter dominated and expansion velocity saturates. Hence for the nonrelativistic equation of state with \( \gamma = 1 \) different scaling laws come out

\[
\Gamma = \text{const}, \quad n \propto r^{-2}, \quad \rho \propto r^{-2}.
\]

Transition between the two regimes (30) and (31) occurs at the radius \( R_c = B^{-1} R_0 \), where \( R_0 \) is initial size of plasma.

Therefore, one may reach the conclusion that for comoving observer the radiation-dominated plasma looks indistinguishable from a portion of radiation-dominated Universe. However, this is true only in the absence of pressure gradients. Strong gradients
are likely present in GRB sources, and they should produce local acceleration in the radiation-dominated electron-positron plasma, making it distinct from the early Universe, where matter inhomogeneities are known to be weak.

It is easy to get from (26) and (28) for internal energy density and temperature

\[
\varepsilon \propto r^{-4}, \quad T \propto r^{-1}, \quad R_0 < r < R_c,
\]

and

\[
\varepsilon \propto r^{-8/3}, \quad T \propto r^{-2/3}, \quad R_c < r < R_{tr},
\]

where \(R_{tr}\) is the radius at which the outflow becomes transparent. The outflow may become transparent for photons also at the acceleration phase, provided that \(R_{tr} < R_c\). For instance, a pure electron-positron plasma gets transparent at the acceleration phase.

3. – Heavy elements

Cosmological nucleosynthesis is a well established branch of cosmology. Classical computations made in the middle of the XXth century revealed that heavy elements cannot be built in the early Universe. Hydrogen and helium contribute approximately 3/4 and 1/4, leaving some room, much less than 1 per cent for deuterium, tritium and lithium. All the heavier elements must have been produced in stars.

Some of these stars, as indicated by observations, end their life as progenitors of GRBs. For this reason it is likely that initially in the source of GRBs elements heavier than hydrogen are present. In this section we consider chemical evolution of plasma in the sources of GRBs.

Assume that in the source of a GRB the amount of energy \(E_0\) is released in the volume with linear size \(R_0\) during the time \(\Delta t\), making this region optically thick to Compton scattering and pair production. The amount of baryons which may be present as well is parametrized by

\[
B \simeq \begin{cases} \frac{Mc^2}{E_0}, & \Delta t \ll R_0/c, \\ \frac{\dot{M}c^2}{L}, & \Delta t \gg R_0/c, \end{cases}
\]

where \(L = dE/dt\) is the luminosity, \(\dot{M} = dM/dt\) is the mass ejection rate and \(M\) is total baryonic mass. Ultrarelativistic outflow is generated through thermal acceleration of baryons by the radiative pressure if plasma is initially energy dominated, i.e.

\[
B \ll 1.
\]

In the case of instant energy release with time interval \(\Delta t \sim R_0/c\) initial temperature in the source of GRB may be estimated neglecting the baryonic contribution, provided
(36) \[ T_0 \simeq \left( \frac{3E_0}{4\pi aR_0^2} \right)^{1/4} \simeq 6.5E_{54}^{1/4}R_8^{-3/4} \text{ MeV}, \]

where \( a = 4\sigma_{SB}/c \), \( \sigma_{SB} \) is the Stefan-Boltzmann constant and the last value is obtained by substituting numerical values for \( E_0 = 10^{54} \text{ erg} \) and \( R_0 = 10^8 \text{ cm} \).

As it has been shown in [7] for temperatures above 1 MeV even low density plasma with density \( n = 10^{18} \text{ cm}^{-3} \) quite quickly destroys all heavier nuclei, and the final state contains just protons and neutrons and some small traces of Deuterium and \( ^4\text{He} \). The timescale of this process (\( \sim 10^{-2} \text{ sec for } T_0 = 1 \text{ MeV} \)) strongly depends on temperature, but the rates of almost all reactions increase with temperature, and correspondingly the abundances of nuclei evolve much faster. Therefore, nuclei disintegration is fast enough to occur before plasma starts to expand and cool on the timescale \( R_0/c \).

During early stages of plasma expansion its temperature decreases in the same way as it happened in the early Universe. Therefore similar synthesis of light elements to BBN occurs also in sources of GRBs. Most important is, however, another similarity with the early Universe: it is well known that practically all free neutrons have been captured into elements heavier than hydrogen. So we do not expect dynamically important free neutrons present in GRB plasma after it started to expand and cool down unless they are engulfed by the expanding plasma later. The role of such free neutrons have been considered in the literature, see e.g. [8] and [9].

4. – Recombination

On the radiation dominated phase both in the early Universe and in GRB plasma entropy conservation (4) results in decrease of temperature. When the comoving temperature decreases below the hydrogen ionization energy, \( E_i = 13.6 \text{ eV} \), the formation of neutral hydrogen is expected.

4.1. Early Universe. – In the early Universe, after the BBN epoch and electron-positron annihilation, cosmological plasma consists of fully ionized hydrogen, helium and small admixture of other light elements. The temperature continues to decrease until it gets sufficiently low to allow formation of neutral atoms: that is the moment in the cosmic history where the formation of the Cosmic Microwave Background Radiation (CMB) happens.

The theory of cosmological recombination of hydrogen, based on three level approximation, has been developed in [10] and [11] in the late 60s. The only modification that such theory undergone in the later years is the account for dark matter and addition of more levels to the model, currently about 300. There is a basic difference with respect to the equilibrium recombination essentially by the process \( e + p \leftrightarrow H + \gamma \), described by the Saha equation

\[ \frac{n_e n_p}{n_H} = \frac{g_e g_p}{g_H} \frac{(2\pi m_e kT)^{3/2}}{\hbar^3} \exp \left( -\frac{E_i}{kT} \right), \]

where \( g_i \) are statistical weights, \( h \) is Planck’s constant. This difference is due to the presence of the 2p quantum level, which produces Ly-\( \alpha \) photons. The absorption of such photons is very strong. However, ionization from the 2p level requires only \( E_i/4 \).
Therefore the formation of neutral hydrogen proceeds through the $2s - 1s$ transition in the presence of abundant Ly-$\alpha$ photons.

In fact, the early Universe would become transparent for radiation even if formation of hydrogen would have been forbidden, see e.g. [12]. The optical depth to Thomson scattering is

$$\tau = \int_{t_0}^{t} \sigma_T n_b c dt \simeq 4 \times 10^{-2} \frac{\Omega_b}{\Omega_m} h \left[ \Omega_{\Lambda} + \Omega_m (1 + z)^3 \right]^{\frac{1}{2}} - 1,$$

(38)

where $\sigma_T$ is the Thomson cross section, $\Omega_i = \rho_i / \rho_c$. $\rho_c = 3H^2c^2/8\pi G$, $H = 100h$ km s/Mpc and $b$, $m$, $\Lambda$ stand for, respectively baryons, dark matter and cosmological constant contributions to the total energy density of the Universe. For large $z$ we have

$$\tau (z_*) = 1, \quad z_* \simeq 8.4\Omega_b^{-2/3} \Omega_m^{1/3} h^{-2/3}.$$

For typical values $\Omega_b h^2 \simeq 0.02$, $\Omega_m \simeq 0.3$, and $h \simeq 0.7$ we have $z_* \simeq 60$. At such redshift the Universe would be expected to become transparent to Thomson scattering. That is exactly what happens in plasma in GRB sources. Below we show that, unlike radiation-dominated cosmological expansion where comoving quantities also fulfill relations (32), the comoving temperature in GRB outflows remains always high enough to prevent recombination of hydrogen.

4.2. GRBs. During both acceleration and coasting phases the comoving temperature decreases with radius, see (32) and (33). The optical depth to Compton scattering may be computed and the corresponding photospheric radius may be obtained, see [13] where ultrarelativistic outflows were analysed in details. In particular, comoving temperature at the photosphere decreases with the baryonic loading $B$ both at acceleration and coasting phases. However, when the outflow reaches the radius $R_s = B^{-2}R_0$, the comoving temperature becomes independent from $B$. In that regime the expression for the photospheric radius is

$$R_{tr} = \left( \frac{\sigma E_0 B}{4\pi m_p c^2} \right)^{1/2},$$

(39)

where $m_p$ and $\sigma$ are proton mass and Thompson cross section, respectively. Indeed, using (36), (39), (32) and (33) in the case of instant energy release we have

$$T_{\text{min}} = BT_0 \left( \frac{R_c}{R_{tr}} \right)^{2/3} =$$

$$\left( \frac{3}{4\pi a} \right)^{1/4} \left( \frac{\sigma}{4\pi m_p c^2} \right)^{-1/3} (E_0 R_0)^{-1/12}.$$

(40)

(41)

Notice how extremely insensitive this value is with respect to the remaining parameters $E_0$ and $R_0$! Expressed in units of typical energy and size

$$T_{\text{min}}^{(s)} \simeq 42 (E_{54} R_8)^{-1/12} \text{ eV}.$$

(42)
In the case of gradual energy release with $\Delta t \gg R_0/c$ and constant luminosity and mass ejection rate the initial temperature is

\[ T_0 \simeq \left( \frac{L}{16\pi\sigma S_B} \right)^{1/4}, \tag{43} \]

and similar expression to (40) may be derived

\[ T_{\text{min}} = \left( \frac{1}{16\pi\sigma S_B} \right)^{1/4} \left( \frac{\sigma}{4\pi m_p c^2} \right)^{-1/3} L^{-1/12} R_0^{1/6} \Delta t^{-1/3}, \tag{44} \]

which may be rewritten, introducing $L = 10^{50} L_{50}$ erg/s and $\Delta t = 1\Delta t_1$ s, as

\[ T_{\text{min}}^{(w)} \simeq 17L^{-1/12} R_0^{5/6} \Delta t_1^{-1/3} \text{ eV}. \tag{45} \]

Even if (45) appears to be less stringent than (42), they are both quite insensitive to initial parameters. As a result, even if the comoving temperature decreases very much compared to its initial value, typically on the order of MeV, at the photospheric radius it is always well above the characteristic temperature 0.3 eV at which recombination happens [14], thus preventing formation of neutral hydrogen. In fact, if such hydrogen would be formed the cross section of interaction of expanding particles with the circumburst medium would drastically decrease. As a consequence no afterglow would be observed.

A simplified way to look at this lower bound on the comoving temperature at the photosphere is to say that if a fraction $\epsilon$ of solar mass is released in the volume having radius $\delta$ solar Schwarzschild radii, then its minimum comoving temperature before transparency is

\[ T_{\text{min}}^{(s)} \simeq 66(\epsilon\delta)^{-1/12} \text{ eV}, \tag{46} \]

in the case of instant energy release and

\[ T_{\text{min}}^{(w)} \simeq 2.8\epsilon^{-1/12} \delta^{1/6} \Delta t_1^{-1/4} \text{ eV}, \tag{47} \]

in the case of gradual energy release during time $\Delta t_1$. Clearly in both cases $\delta > 1$, and likely $\epsilon < 1$. Notice, that while in the case of instant energy release the lower bound on temperature decreases with increasing $\delta$, it instead increases in the case of gradual energy release.

The baryon to photon ratio in GRB plasma like in cosmology is large. This ratio may be estimated as

\[ \frac{n_\gamma}{n_B} = \frac{m_p}{m_e} \frac{1}{B \langle \varepsilon \rangle} \simeq 1.8 \times 10^5 B_{-2}^{-1/2} \langle \varepsilon \rangle^{-1}, \tag{48} \]

where $B_{-2} = 10^{-2} B$, and $\langle \varepsilon \rangle$ is average photon energy in the source of GRB in units of electron rest mass energy. Thus the optical depth of electrons is much larger than the one of photons and it is given in [15]

\[ \frac{\tau_e}{\tau_\gamma} = \log \Lambda + \frac{n_\gamma}{n_e} \approx \frac{n_\gamma}{n_B}, \tag{49} \]
where $\Lambda$ is the Coulomb logarithm. It means that electrons are kept in equilibrium with photons when the latter already decoupled from them [16]. In other words, electrons are forced to keep the local temperature of photons. This may lead to efficient Comptonization of the photon flow when it is decoupled from plasma and is passing through electrons having locally different temperature.

As soon as plasma gets collisionless, laboratory spectrum of photons, baryons and electrons is maintained. If it was thermal at decoupling it will remain so. This shows another difference with respect to cosmology, where energy of all particles decoupled from the thermal bath decreases due to the cosmological expansion. For that reason in cosmology only the shape of the spectrum is conserved with expansion, but not the temperature.

Therefore, we have reached the conclusion that hydrogen recombination which is responsible for transparency of cosmological plasma does not occur in GRB plasma. This difference in physical conditions may result in deviations from black body spectrum, as observed in GRBs. Available studies of photospheric emission in GRBs in the literature show that deviations from the perfectly thermal spectrum come mostly from three effects: a) dynamical and ultrarelativistic character of plasma outflows and geometric effects [13]; b) "fuzzy photosphere" effect [17, 18] and c) possible dissipation mechanisms at the photosphere [9, 19, 20].

Recently we presented a theory of photospheric emission from relativistic outflows, see [13]. Assuming that the spectrum of radiation in the comoving reference frame is the perfect black body one, we have shown that the spectrum seen by a distant observer may be essentially nonthermal due to both geometric and dynamical special relativistic effects. The possibility that the spectrum of photospheric emission is nonthermal also in the comoving frame is under investigation.

5. – Conclusions

Regarding the dynamical aspects, there is an apparent similarity between the electron-positron plasma in the early Universe and the one in GRB sources. For an observer comoving with the radiation-dominated plasma in GRB source it may look indistinguishable from a portion of radiation-dominated Universe. However, this is true only in the absence of pressure gradients. Strong gradients are likely present in GRB sources, and they should produce local acceleration in the radiation-dominated electron-positron plasma, making it distinct from the early Universe, where matter inhomogeneities are known to be weak.

There is also an apparent similarity with respect to the nucleosynthesis phenomenon. Given that the temperature reached in GRB sources, see Eq. (36), may be as high as several MeV, nuclear reactions are expected to operate on timescales of $10^{-2}$ sec or shorter. That is on the order of magnitude of dynamical timescale of the GRB sources. It means that reprocessing of nuclear abundances may likely take place in GRB sources. Since observations imply that GRBs may originate from compact stellar objects elements heavier than helium are likely to be present in GRB sources. Such heavy elements are then destroyed, resulting mainly in protons with small admixture of helium. Thus, similarly to the early Universe, we do not expect dynamically important free neutrons present in GRB plasma after it started to expand and cool down unless they are engulfed by the expanding plasma later.

Finally, there is an important difference between the electron-positron plasma in the early Universe and the one in GRB sources. We show in this paper that unlike the
primordial plasma which recombines to form neutral hydrogen, and emits the Cosmic Microwave Background Radiation, GRB plasma does not cool down enough to recombine. Therefore GRB plasma always becomes transparent due to Compton scattering.

REFERENCES

[1] R. Ruffini, G. Vereshchagin, and S.-S. Xue, Physics Reports 487, 1 (2010).
[2] A. G. Aksenov, R. Ruffini, and G. V. Vereshchagin, Phys. Rev. Lett. 99, 125003 (2007).
[3] S. Weinberg, *Cosmology*, Oxford University Press, April 2008., 2008.
[4] T. Piran, A. Shemi, and R. Narayan, MNRAS 263, 861 (1993).
[5] R. Ruffini, J. D. Salmonson, J. R. Wilson, and S.-S. Xue, A&A 359, 855 (2000).
[6] A. Shemi and T. Piran, ApJ 365, L55 (1990).
[7] E. Kafexhiu, Excitation and destruction of nuclei in hot astrophysical plasmas around black holes, in *25th Texas Symposium on Relativistic Astrophysics*, 2010.
[8] E. V. Derishev, V. V. Kocharovsky, and V. V. Kocharovsky, Nuclear Physics B Proceedings Supplements 80, C612+ (2000).
[9] A. M. Beloborodov, MNRAS 407, 1033 (2010).
[10] Y. B. Zeldovich, V. G. Kurt, and R. A. Syunyaev, Zhurnal Eksperimental no i Teoreticheskoi Fiziki 55, 278 (1968).
[11] P. J. E. Peebles, ApJ 153, 1 (1968).
[12] P. D. Naselsky, D. I. Novikov, and I. D. Novikov, *The Physics of the Cosmic Microwave Background*, Cambridge Astrophysics, Cambridge Univ. Press, 2011.
[13] R. Ruffini, I. A. Siutsou, and G. V. Vereshchagin, arXiv:1110.0407 (2011).
[14] J. Bernstein and S. Dodelson, Phys. Rev. D41, 354 (1990).
[15] S. R. de Groot, W. A. van Leeuwen, and C. G. van Weert, *Relativistic kinetic theory. Principles and Applications*, North Holland Publishing Company, 1980.
[16] O. M. Grimsrud and I. Wasserman, MNRAS 300, 1158 (1998).
[17] A. Pe'er, ApJ 682, 463 (2008).
[18] A. M. Beloborodov, ApJ 737, 68 (2011).
[19] K. Toma, X.-F. Wu, and P. Mészáros, MNRAS 415, 1663 (2011).
[20] F. Ryde et al., MNRAS 415, 3693 (2011).