GRBs from unstable Poynting dominated outflows

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**ABSTRACT**

Poynting flux driven outflows from magnetized rotators are a plausible explanation for gamma-ray burst engines. We suggest a new possibility for how such outflows might transfer energy into radiating particles. We argue that the Poynting flux drives non-linearly unstable large amplitude electromagnetic waves (LAEMW) which “break” at radii \( r_t \sim 10^{14} \text{ cm} \) where the MHD approximation becomes inapplicable. In the “foaming” (relativisticly reconnecting) regions formed during the wave breaks the random electric fields stochastically accelerate particles to ultrarelativistic energies which then radiate in turbulent electromagnetic fields. The typical energy of the emitted photons is a fraction of the fundamental Compton energy \( \epsilon \sim f \hbar c/r_e \) with \( f \sim 10^{-3} \) plus additional boosting due to the bulk motion of the medium. The emission properties are similar to synchrotron radiation, with a typical cooling time \( \sim 10^{-4} \text{ sec} \). During the wave break, the plasma is also bulk accelerated in the outward radial direction and at larger radii can produce afterglows due to the interactions with external medium. The near equipartition fields required by afterglow models maybe due to magnetic field regeneration in the outflowing plasma (similarly to the field generation by LAEMW of laser-plasma interactions) and mixing with the upstream plasma.
1. Introduction

A Poynting flux driven outflow from a magnetized rotator is a promising paradigm for gamma-ray burst (GRB) engines and there have been various implementations of this concept (c.f. Usov 1992; Thompson 1994; Blackman et al. 1996; Mészáros & Rees 1997; Kluzniak & Ruderman 1998). Such models are appealing in light of the fact that neutrino driven GRB emission from compact object mergers (Ruffert & Janka 1998; Janka et al. 1999) might fall short in radiative luminosity. Also, Poynting flux models may provide a source of magnetic fields which could help alleviate the field amplification problem in GRBs themselves and in afterglows which requires fields of the order of equipartition (Wijers & Galama 1998; Frail et al. 1999; Granot et al. 1999). Scenarios that could produce Poynting flux dominated outflows (PFDOs) require a source of magnetic fields \( \Omega \gtrsim 10^{15} \text{ Gauss} \), and rotation speeds of order \( \Omega \sim 10^{4} \text{/sec} \). The total available energy from a compact \((R_{0} \sim 10^{6} \text{ cm}, M \gtrsim M_{\odot})\) object is then \( M \Omega^{2} R_{0}^{2}/2 \gtrsim 10^{53} \text{ ergs} \) and a dipole type luminosity \( 2B_{0}^{2} R_{0}^{6} \Omega^{4}/3c^{3} \sim 2 \times 10^{50} \text{ ergs/sec} \). Such combinations could be generated in a number of plausible cases: an accretion torus surrounding a black hole that formed from a neutron star merger (Rees & Mészáros 1997), strongly magnetized neutron stars (“magnetars” Duncan & Thompson 1992) possibly formed from accretion induced collapse of a white dwarf with possible dynamo amplified fields in the hot young neutron star (Usom 1992; Duncan & Thompson 1992, Blackman et al. 1996), failed supernova Ib (Woosley 1993), rapidly rotating magnetized black holes undergoing a Blandford-Znajek type energy extraction from supernovae (Lee et al. 1999, Brown et al., 2000) or hyper-accreting accretion disks (Popham et al. 1999), with an MHD dynamo (Araya-Gochez 2000).

Without further specifying the nature of the central engine we suppose that GRB are due to such stellar mass objects, so that the conditions required for the generation of PFDOs are satisfied. Here we are interested in how Poynting flux models accelerate particles and produce gamma-ray emission.

The suggestion that GRBs are PFDOs whose energy is converted to particles and escaping radiation only at large distances from the central rotator potentially accounts for some important characteristics of GRBs (cf. Blackman et al. 1996, Rees & Mészáros 1997) (i) mass loading - PFDO can be launched without much matter, and particularly without much baryon contamination (as in pulsars, the wind may carry no baryons at all); (ii) compactness problem - the PFDO converts its energy effectively into particles and gamma-rays at distances many orders of magnitude from the central engine, (iii) collimation - PFDO have a preferred direction along the rotation axis of the central engine; (iv) jets observed in pulsars and (probably) AGNs provide examples of operation of PFDO; (v) the energy from Poynting flux can be easily converted into high frequency electromagnetic radiation. The presence of strong magnetic fields from the rotator provides electromagnetic energy which might be tapped or converted into the fields which accelerate the radiating particles.

PFDOs may differ from the conventional blast wave models in several important respects. The lepto-baryonic content of the outflow in the inner region is expected to be small by analogy to pulsar winds, and the energy carried by matter in the central region will be less than the energy carried by EM fields. As in the relativistic pulsar wind, the only other well established PFDO, the baryons may be absent altogether while energy flux in the electron-positron pairs may be as small as one millionth of the Poynting flux.

The conventional estimates of the relativistic expansion Lorentz factors and density of the outflowing plasma used in the framework of blast wave models may not be applicable to the Poynting dominated outflows because of the strong magnetic fields (Paczynski 1990, Usov 1999). Strong fields dramatically reduce the importance of pair production from photon-photon interactions relative to the photon-magnetic field interactions. Other exotic QED processes, like photon splitting (e.g., Baring 1991, Adler 1971), will also become important for magnetic fields larger than the critical value of \( 4.4^{13} \text{ Gauss} \). At this moment, we are unaware of the calculations of the Compton scattering cross section in supercritical magnetic fields, but we expect that it will also be strongly suppressed.

The above suggests that in the central regions of PFDOs, the magnetic pair production and photon splitting may be more important than pair production in photon collisions. Note that this situation is realized in pulsars. There acceleration of particles to energies beyond \( 10^{6} \text{ MeV} \) proceeds relatively quietly, without producing a large number of observable photons. Very few rotationally powered pulsars are observed at high energies. Thus conventional estimates
of the Lorentz factors and densities in GRBs based on two photon pair production may be irrelevant.

To estimate a “revised” density of the generated pairs in the strong magnetic rotator case applicable to GRB, we use the analogy with pulsars where the number of pairs produced is related to the local Goldreich-Julian density \( n_{GJ} = \Omega B/(2\pi \kappa c) \) (Goldreich & Julian 1969). The pair number density is usually \( \lambda \sim 10^3 - 10^5 \) times larger than \( n_{GJ} \) in the inner regions. The factor \( \lambda \) depends sensitively on the magnetic field, its curvature and accelerating potential. A highly curved magnetic field produces denser and less relativistic outflows.

In this paper we address the fundamental but unresolved issue of how Poynting flux is converted into gamma-rays and the subsequent emission characteristics. In section 2 we give the basic picture of the magnetized rotator magnetosphere and the propagation of large amplitude electromagnetic waves. In section 3 we discuss the breaking of these waves, and the predicted emission from random electromagnetic fields is discussed in section 4. There we address the spectrum and time dependence of the emission. The afterglow possibilities from our picture are addressed in section 5. We conclude in section 6.

2. Basic Picture

Here we discuss the basic proposed scenario for the evolution of PFDOs. If there is a strong overall magnetic dipole component, then near the central engine the PFDO will resemble a pulsar. Along the magnetic poles it will produce jet-like flow with helical magnetic field, oscillating on scale length of \( \sim 10^3 \) cm. Initially Poynting energy flux dominates over the particle flux by a ratio \( \sim 10^6 \) typical for pulsars.

Inside the “light cylinder” \( r_{LC} = (c/\Omega) = 3 \times 10^6 \) and the magnetic field falls off as \( r^{-3} \), while beyond the light cylinder the magnetic field is dominated by the toroidal component which falls of as \( r^{-1} \). Thus at large distances \( r \gg r_{LC} \) we have

\[
B_\phi = B_0 \frac{R_0}{r} \left( \frac{\Omega R_0}{c} \right)^2
\]

where \( R_0 = 10^6 \) cm and \( B_0 = 10^{15} \) G are the size of the central engine and the initial magnetic field.

Following the analogy to pulsar winds (e.g. Usov 1992; 1994), at small distances the magnetic field of the wind is frozen in the plasma. As the plasma flows out, the plasma density decreases in proportion to \( r^{-2} \) reaching a radius \( r_t \), where it becomes less than the critical charge density \( n_t \) required for applicability of MHD (Goldreich & Julian 1969; Michel 1969, 1994; Coroniti 1990; Melatos & Melrose 1996). Locally the critical charge density \( n_t \) is equal to the Goldreich-Julian density \( n_{GJ} \). The latter decreases only as \( r^{-1} \).

In what follows, we normalize the real density of the PFDO \( n \) to the Goldreich-Julian density at the light cylinder \( n_{GJ}(r_{LC}) \), that is

\[
n = \lambda n_{GJ} = \frac{\Omega B_0}{2 \pi \kappa c} \left( \frac{R_0 \Omega}{c} \right)^3 = 4 \times 10^{15} \text{cm}^{-3}\]  

Since the plasma density satisfies \( n(r) = n(r_{LC})(r_{LC}/r)^2 \), while \( n_{GJ}(r) = n_{GJ}(r_{LC})(r_{LC}/r) \), the plasma density equals Goldreich-Julian density at

\[
r_t = \lambda r_{LC} = 10^{14} \text{cm},
\]

where we have chosen \( \lambda = 3 \times 10^7 \) (see also Usov 1994; Blackman & Yi 1998). The density and magnetic field at the transition radius are

\[
n_t = n_{GJ}(r_t) = 10^8 \text{cm}^{-3}, \quad B = 10^6 \text{G}.
\]

The physical parameters near the transition region are listed in Table I.

At \( r > r_t \) the the wind field is transformed into a Large Amplitude Electromagnetic Wave (LAEMW). The dimensionless strength parameter of the LAEMW, \( \nu_0 = \epsilon E/(mc\Omega) \), near the transition point is

\[
\nu_0 = \frac{\Omega R_0}{c^2} \frac{B_0}{n} \sim 10^9.
\]

Table 1: Physical parameters near the transition region

| Parameter | Value |
|-----------|-------|
| \( B, \text{ Gauss} \) | \( 10^6 \) |
| \( \omega_B, \text{ rad/sec} \) | \( 2 \times 10^{13} \) |
| \( \nu_0 = \omega_B/\Omega \) | \( 2 \times 10^9 \) |
| \( r_L = \frac{c}{\omega_B}, \text{ cm} \) | \( 1.5 \times 10^{-3} \) |
| \( \omega_p = \sqrt{2\Omega \omega_B}, \text{ rad/sec} \) | \( 6 \times 10^8 \) |
| \( r_s = \frac{c}{\omega_p}, \text{ cm} \) | \( 50 \) |
| \( r_{ts} = \frac{\omega_B}{\omega_p} = \sqrt{\frac{2}{\nu_0}} \) | \( 3 \times 10^{-5} \) |
| \( \frac{\nu_0}{\nu_0} \) | \( 5 \times 10^9 \) |
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lates an electromagnetic shock. Interaction of such

active medium (Agrawal 1989). A wave overturn cre-

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of the foam). As the accelerated relativistic particles

trum of random EM fields. It is an electro-

We describe this instability in the next section.

Once the EM wave overturns, it creates a broad

spectrum of random EM fields. It is an electro-
magnetic analog of the foam of the deep water sur-

face waves. Overturn of the initial coherent wave

will create random EM fields and currents which will

be strongly dissipated. In general, there will be a

nonzero component of electric field along magnetic

field. We will call this state a relativistically strong

electro-magnetic turbulence. Since the end result is

the dissipation of the magnetic energy, this regions

may also be called a relativistically strong reconnect-
ing cites. Particles will be accelerated in this EM

turbulence to relativistic energies (in the rest frame

of the foam). As the accelerated relativistic particles

move through the EM foam they generate the ob-

served GRB emission. The emission is due to the ac-

celeration in the random E-B fields. Emission would

emanate from multiple breaking regions. The sepa-

rate spikes in the GRB profiles would be due to sepa-

rate "foaming regions" - separate regions of the wave

break.

In many respects the "foaming regions" resemble

optical shocks observed in the laboratory when an

optical pulse steepens as it propagates in an optically

active medium (Agrawal 1989). A wave overturn cre-

ates an electromagnetic shock. Interaction of such

EM shocks will unavoidably lead to a creation of inter-

nal plasma discontinuities with strong dissipative ef-

fects. The "foaming regions" play the role of “internal

shocks” (c.f. Piran 1999 for a review) of GRB theory.

This term is usually used to mean “internal processes

in the blastwave which generate bursts” not necessarily

a specific mechanism, which we are providing here.

Alternatively, the "foaming regions" maybe thought of

as relativistically reconnecting regions, where the en-

ergy of magnetic field is converted into particles in a

relativistically turbulent manner.

It is often assumed that some kind of Fermi scat-

tering process operates at the GRB relativistic shocks.

In the shock Fermi process, particle scattering occurs

across the shock between “scattering centers” of col-

lections of e.g. Alfven waves. Technically speaking,

shock Fermi acceleration is not the likely mechanism

for particle acceleration for generic relativistic shocks

because for upstream magnetic field angles greater

than 1/Γ to the shock normal, where Γ is the bulk

upstream Lorentz factor, the particles would have to

travel faster than c to get back upstream along field

lines. Thus unless the shock is strictly parallel (un-

likely situation) then shock Fermi acceleration won’t

work.

Hoshino et al. (1992) studied 1-D models of

positron acceleration in relativistic perpendicular shocks

(where the field is perpendicular to the shock normal)

in ion-electron-positron plasmas. Their mechanism

involves the absorption, by positrons, of gyro-phase

bunched magnetosonic waves emitted by protons re-

lected at the shock front. For ~ 10% ion fraction,

non-thermal tails are seen in positrons. They dis-

cuss limitations of their simulations with respect to

their choice of electron-position number ratio, the fact

that the simulations are 1-dimensional, and that they

do not consider oblique shocks (shocks for which the

fields are not exactly perpendicular to the shock nor-

mal). They point out that electric fields may play

more of a role in particle acceleration in the latter

case. In general, they agree that more work is needed.

We entertain the possibility that electric fields may

be important for acceleration, but not necessarily only

across a thin shock. Rather, we suggest that as an

LAEMW overturns, the acceleration, even in pair

plasmas, might be done directly by turbulent elec-

tric fields (see also Lembege & Dawson 1987 for over-

turn discussion). The thickness of the turbulent re-

gion may be much larger than the thickness of a typi-

cal shock. The overturn may force the region to be

highly strongly turbulent over many gyro-radii.
The picture of the "foaming regions" bears some similarity to the "plasma turbulent reactor" (Tsy-tovich & Chikhalev, Norman & ter Haar 1975; Rees 1967). In both cases there is a strong interaction of plasma kinetic turbulence with electromagnetic waves which may lead to a formation of stationary distribution of particles. But the "plasma turbulent reactor" has been studied only in the limit of weak turbulence theory, when the energy density of the turbulent plasma fluctuations is much smaller than the total energy density of plasma and when plasma is optically thick to transverse EM waves. The assumption of weak turbulence allowed for a closed self-consistent solution for the plasma turbulent spectrum and particle distribution. On the other hand, the turbulence in the foaming regions of GRBs should be very strong, with an energy density much larger than that of the plasma. Detailed properties of such superstrong turbulence such as the stationary spectra, are not known. Under certain conditions we will show that the emission from such turbulence does not depend on its detailed properties (see section 4). We will not restrict ourselves to a particular model of the turbulence, but we do discuss Phillips-type and relativistic hydrodynamics-type spectra in section 4.4 and 4.5.

Each overturning region produces a burst of gammarays. The temporal behavior of such a burst might be most similar to a FRED (fast rise exponential decay) curve, since the fraction of LAEMW energy going into turbulence and particle acceleration likely occurs more rapidly than the subsequent radiation drain. The overall duration of a microburst is of the order of micro seconds (Eq. (24)). To account for the energetics (\( \epsilon_{mb} \approx 10^{48} \text{ergs per microburst} \) with of order \( 10^4 \) microbursts as discussed later in section 4) the overturning region should draw energy from a volume \( V \sim 10^{16} \text{cm}^3 \), i.e. the typical size of the turbulent region should be \( \sim 10^{12} \text{ cm} \) which is much smaller than the distance to the central region. Two other arguments may influence our estimates of the emitting region: (i) using the analogy with surface foam we can expect the transverse dimensions of the foaming region to be larger than the longitudinal dimension, (ii) during a break of a wave packet energy from many wave periods is "piled up" in a narrow breaking region, reaching the thickness of one wavelength \( \sim 10^6 \text{ cm} \). This value also turns out to be of the order of \( c \) times the radiative decay length of a typical particle.

The total number of overturning regions will depend on the structure of the flow, which is in turn determined by the structure of the central engine. For example, if the PFDO has a large dipole component, then the structure of the flow will be more regular and we expect only several regions of wave overturn. On the other hand if the structure of the magnetic field near the central engine is more multipolar, we expect many "foaming regions".

At the transition radius, momentum conservation dictates that there will remain some energy in the LAEMW which can be used for bulk plasma acceleration. The bulk motion of the outflowing plasma should drive a relativistic shock wave into the ambient gas. Synchrotron emission of particles at this shock could produce the observed afterglows. This would coincide with the conventional external shock models of afterglows (c.f. Piran 1999), while providing a possible improvement: as we discuss further in Section 5, the outflowing plasma beyond the transition radius may still be strongly magnetized, thus supplying the elusive magnetic field required in the synchrotron models of afterglows. In this case the afterglows are something like the analogs of the wisps observed around the Crab pulsar (the wisps are thought to be the location of the reverse shock where the pulsar winds pressure is balanced by the pressure of the supernova ejecta; Gallant & Arons 1994). Alternatively, the afterglows may be generated in a process similar to the one considered by Smolsky & Usov (2000), when the particles from the external medium are reflected from the strong magnetic field of the ejecta. In both cases however, we point out that the outflow can supply the magnetic field to the afterglow emitting region.

3. Overturn of LAEMWs

Here we consider the transition of the PFDO from MHD to the wave regime. In the MHD approximation, the velocity of the wind equals the velocity of the field patterns. This wind velocity can be described as a wave phase velocity with trapped particles. In the laboratory frame this velocity is determined by the electric drift in the crossed electric and magnetic fields, \( v_d = \frac{E \times B}{B} = \frac{E}{B} < 1 \). For the PFDOs, the leading EM terms are \( B_\phi, E_\theta \propto \frac{1}{r} \) so that the drift velocity (the velocity of the outflow) is directed along

\[ \frac{E}{B} < 1 \]
the radial direction $r$. When the MHD approximation breaks down, the phase velocity (i.e. the velocity of the field pattern) will grow to $v_{ph} = \frac{c}{\beta} > 1$. Thus, during the transition from MHD approximation to plasma wave solution the phase velocity of the field patterns will increase from subluminal to superluminal. For plasma particles this process will look as if the EM wave is coming into plasma "from minus infinity". As the EM wave penetrates into the plasma it exerts ponderomotive force along the direction of its propagation which provides an acceleration for the plasma. Below we show that this process of acceleration is unstable toward wave overturn.

The equations describing the evolution of the LAEMW waves in plasma are Maxwell’s equations and the mass and momentum continuity equations for electron and positron fluids. To show the wave overturn simply, we consider a cold plasma. This is a reasonable approximation when the particles acquire relativistic motion and momentum continuity equations for electron and $p_{\perp e,p} = \pm e a_\perp. \quad p_z = \text{const} \quad (10)$

Thus the transverse velocity of the plasma electrons and positrons are antiparallel and directed along the instantaneous magnetic field, but perpendicular to the electric field of the wave. The longitudinal velocity is arbitrary.

To first order in the small parameters we find
\begin{align}
a_z = 0, \quad \phi = 0 \\
\frac{\partial_t p_z + \partial_z \gamma}{\gamma} = 0, \quad \gamma = \sqrt{1 + \nu^2 + p_z^2} \quad (11)\\n\frac{\partial_t n + \partial_z (n v_z)}{v_z} = 0 \quad (12)
\end{align}

and expand the dynamic equations in small quantities $\frac{\partial_a a_n}{\omega a_n}$ and $\frac{\partial_a a_n}{\nu a_n}$. In nonlinear optics this is known as a slowly varying envelope approximation (e.g. Agrawal 1989). To further simplify the treatment, we consider a circularly polarized EM wave for which the energy of particles in a wave remains constant during oscillations (polarm outflows from magnetized rotator should produce circularly polarized EM waves). For linearly polarized waves, an additional averaging over one period of oscillation is needed. In the zeroth order, we find

\begin{equation}
0 = \left( e a_\perp + e A \right) - \frac{\partial_x p}{\gamma} = \frac{\partial_x p}{\gamma} \Perp (9)
\end{equation}

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\end{align}

and a Schrödinger type equation for the transverse vector potential (e.g. Berezhiani et al. 1992). Note that in the unmagnetized pair plasma, the variation of amplitude along the direction of wave propagation does not create charge separation, so $n \equiv n_{e,p}$ and $p_z \equiv p_{z e,p}$.

Equations (1) and (2) should be integrated along the group velocity characteristics $\partial_z / \partial t = v_g$. Assuming that all quantities depend on the variable $\xi = z - v_g t$, equations (1) and (2) become
\begin{align}
- v_g \partial_t p_z + \partial_z \gamma = 0 \quad (13)\\n\partial_t \ln n = \frac{\partial_z v_g}{v_g} \qquad (14)
\end{align}

For nonlinear waves, the group velocity $v_g$ depends on the Lorentz factor of the particle flow and on the local plasma density, $v_g(n, \gamma)$. The general solution of the system (13-14) for varying $v_g(n, \gamma)$ seems to be untractable, so we first analytically illustrate the overturn for the case of constant group velocity and later resort to numerical integration for the fully nonlinear case.

Integrating (13) along the trajectories $v_g = \text{const},$
we find (Clemmow 1974)
\[
v_z = \frac{n_g \left( 1 + n_0^2 + \sqrt{1 - (-1 + n_g^2) \nu_0^2} \right)}{1 + n_g^2 + \nu_0^2}
\]
\[
\gamma = \frac{n_g^2 + \sqrt{1 - (-1 + n_g^2) \nu_0^2}}{n_g^2 - 1}
\]
where \( n_g = 1/\nu_g \) and \( \nu_0 \) is the dimensionless amplitude of the wave. Equation (15) shows that for the intensity of the wave above the threshold value
\[
\nu_0 = \nu_g/\sqrt{1 - \nu_g^2} \quad (\text{at which point} \quad v_z = v_0 = \frac{\nu_0}{\sqrt{1 + \nu_0^2}})
\]
the Lorentz factor \( \gamma \) becomes undetermined. This implies that hydrodynamic description of the medium is no longer valid: the waves overturn (Dawson 1959, Lembege & Dawson 1987; Lembege & Dawson 1989). During wave overturn, a mixing of various parts of the wave occurs, which destroys its oscillatory structure. Note that in the overturning regions the singularity in the density (14) at this approximation is marks the regime of strong dissipation, where the bulk motion is randomized. It is thus a regime of particle acceleration.

To further illustrate the wave overturn, we numerically integrate (11) and (12) for a given profile of the incoming wave (Fig. 1).

The overturn of strong nonlinear waves is a well known phenomena in fiber optics where optical shocks (sometimes called kink solutions) form when a narrow intense laser pulse propagating through a nonlinear medium (e.g. Agrawal 1989; Rothenberg & Grishkowsky 1989; Trippenbach & Band 1998). Unlike relativistically strong LAEMW, the typical nonlinear intensities of the short laser pulses are small and usually a nonlinear Schrödinger equation is used to consider wave breaking (Agrawal & Headly 1992).

In an electron-ion plasma, the ponderomotive force felt by ions and electrons is different, resulting in different accelerations and the creation of large unstable currents. Generally, electron-positron plasmas are more stable to parametric and modulational instabilities than ion electron plasmas since some of the nonlinear 3rd order currents cancel out. However for a magnetized pair plasma the ponderomotive force acting on electrons and positrons is different. Thus if the outflowing plasma of GRBs has a large ion component or large nonoscillating magnetic fields, then the propagation of LAEMWs will be significantly more unstable.

4. Emission in a random EM field

4.1. General remarks

In this section we consider generation of X-ray and gamma-ray emission from the acceleration and propagation of relativistic electrons/positrons through relativistically strong electromagnetic turbulence. Acceleration of particles is due to random alignments of the electric field along the particle’s velocity. Radiation, on the other hand, is due to the random components of the magnetic and electric field perpendicular to the particle’s velocity.

The typical electric field \( E \sim B \) produced near the transition radius is orders of magnitude larger than the Dreicer field in the mildly relativistic plasma with critical density \( n_{GJ} \): \( E_D \sim \frac{\nu_g}{c} \sim \left( \frac{\nu_g}{c} \right) \). For electric fields larger than Dreicer field, Coulomb collisions between electrons can be neglected and they are freely accelerated by the field.

The spectrum of the plasma particles will be determined by the competition between acceleration and emission in the stochastic EM fields. If these two process balance, a stationary spectrum will be reached. In this case the energy of particles would be radiation limited and only weakly dependent on the details of the acceleration mechanism.

Acceleration of particles may proceed in two regimes which we refer to as “maximally efficient” and “stochastic acceleration” respectively. When the typical scale of the field fluctuations is larger than the typical radiative loss length, then the particle acceleration can be thought of as acceleration in an almost constant DC electric field and thus will be maximally efficient. If, on the other hand, the typical radiative length is much larger than the coherence length of the field fluctuation, then the particle will encounter multiple electric field fluctuations and reversals before attaining steady state energy. Here the acceleration is stochastic. We expect that electromagnetic turbulence to be broadband, extending over scales from the skin depth up to the size of the turbulent region. The relative importance of the maximally efficient and stochastic acceleration types depends on the (unknown) spectrum of the strongly relativistic turbulence.

Emission from a relativistic particle in random \( E \) and \( B \) fields may be qualitatively related to emission from an electron in a random \( B \) field alone. To see this we note that the acceleration due to the electric field along the particle’s velocity \( E \parallel v \) is much less
(by a factor $\sim \gamma^2$) than the acceleration due to transverse component of the electric field $E_L$. (Landau & Lifshitz 1975). The component of the electric field $E_L$ transverse to the particle’s velocity acts similarly to a $B$ field. Thus the emission of an electron in random $E$ field will be very similar to the emission of an electron in effective magnetic field of the same magnitude. Moreover, in a relativistic plasma the random electric and magnetic fields should be of the same order.

The particle emission properties in turbulent magnetic and electric fields also depend on the typical scales of field fluctuations. Relativistically moving particles emit most of the radiation along the particle’s velocity within an angle $\delta \theta \sim 1/\gamma$. The resulting spectrum of the emission depends on whether the total deviation angle of the particle during its transition through EM field is larger or smaller that $\sim 1/\gamma$. For transverse acceleration, the relevant scale is the relativistic Larmor radius $r_L = c\gamma/\omega_{eff}$. Perturbations of the accelerating fields with scales larger than $r_L$ produce deflection angles larger than $\sim 1/\gamma$. The resulting emission will be similar to synchrotron emission. The total spectrum from an emitting plasma is then the average over the emitting volume of synchrotron spectra.

On the other hand, for field fluctuations with wavelengths smaller than the Larmor radius the deflection during emission is smaller than $\sim 1/\gamma$. The particle moves almost straight along the line and experiences high-frequency jitter in the perpendicular direction from the random Lorentz force. Thus produces “jitter” radiation (Medvedev 2000, Landau & Lifshitz II) with a typical frequency $\sim \gamma^2 c/\lambda$, where $\lambda$ is the scale length of perturbations. The spectrum is then determined by random accelerations of the particle. As we show in Section 4.3 however, the expected turbulent scales are much larger than the Larmor radii, so that the emission will resemble synchrotron, not jitter, emission.

### 4.2. Order of magnitude estimates

To illustrate the some of the main points discussed above, we first consider the case when the smallest scale of the field fluctuations is assumed to be larger than the radiative loss length, which in turn is assumed to be larger than the non-relativistic Larmor radius. In this case, the acceleration and emission result from quasistatic electric and magnetic fields.

Also in this case, the typical frequency and energy loss of a particle follows from the analogy to synchrotron radiation in the an effective magnetic field $B_{eff} = \sqrt{< B^2 > + < E^2 >}$, where $< B^2 >$ and $< E^2 >$ are the rms of the fields:

$$\omega_t \sim \gamma^2 \omega_{eff},$$

and where $\omega_{eff} = c\sqrt{< B^2 > + < E^2 >}/mc$ is the "effective cyclotron frequency". The total emissivity from a particle is

$$\frac{dU}{dt} \sim \gamma^2 \frac{e^2}{c} \omega_{eff}^2.$$  

In a steady state, the radiative losses would be exactly compensated by acceleration in the fluctuating electromagnetic fields. Assuming maximally effective acceleration, we find the energy gain to be

$$\frac{dU}{dt} \sim \frac{1}{\sqrt{2}} mc^2 \omega_{eff},$$

where the factor $1/\sqrt{2}$ is due to the fact that only the electric field contributes to acceleration. Equating Eqs. (17) and (18) we find the typical Lorentz factor of the particles to be

$$\gamma \sim \sqrt{\frac{r_L}{2r_e}} \sim 7 \times 10^4,$$

where $r_L = c/\omega_{eff}$ is the effective Larmor radius and $r_e = e^2/mc^2$ is the classical electron radius. The typical emission energy is then given by a combination of fundamental constants only:

$$\epsilon \sim \frac{1}{\sqrt{2}} \frac{h c}{2r_e} \sim 40\text{MeV}. $$

Thus the typical energy of emission of a relativistic particle in turbulent EM field turns out to be independent of the spectral details of the turbulence in the maximally efficient regime. The energy of the particles is radiation limited.

To see this more clearly, we compare the energy density in the particles in a stationary state with the equipartition energy density that would be achieved if approximately half of the energy of the initial magnetic field were transformed into particle energy. The ratio of the magnetic to particle energy density may be expressed as a ratio of the relativistic Larmor and skin depths

$$\eta \equiv \frac{U_m}{U_p} = \left( \frac{\rho_s}{\rho_L} \right)^2 = \frac{1}{< \gamma >} \left( \frac{r_s}{r_L} \right)^2 >> 1, $$

where $\rho_s$ and $\rho_L$ are the skin depths due to the magnetic and electric fields, respectively.
where \( r_s^{(r)} = \sqrt{< \gamma > r_s} \) and \( r_L^{(r)} = \gamma r_L \) are the relativistic skin depth and the relativistic Larmor radius and the latter relation follows for the relevant parameter choices. Since both quantities are properties of the plasma as a whole, we used the averaged Lorentz factor of the particles \( < \gamma > \) in their definitions. Had we assumed equipartition, \( \eta \sim 1 \), the average energy \( < \gamma > \) would have been

\[
< \gamma > \sim \left( \frac{r_s}{r_L} \right)^2 = 10^9
\]  

(22)

which is much larger than \( 10^3 \). Thus the typical energy density in particles is radiation limited and is less than the energy density in EM fields by a factor \( \eta \sim 10^4 \).

The typical synchrotron decay time for a single particle is very short

\[
t_r \sim \frac{c}{r_e \gamma \omega_{ef}^2} \approx 5 \times 10^{-8} \text{sec}
\]  

(23)

Yet in a steady state the energy lost by a particle due to radiation is resupplied by the energy stored in turbulent EM fields. The typical radiation decay times for the turbulence then will be \( \eta \) times longer

\[
\tau_r \sim \eta t_r = 5 \times 10^{-4} \text{sec}
\]  

(24)

This radiation decay time maybe related to the duration of micro spikes/bursts - each micro spike being due to one overturning region.

4.3. Short turbulence scales: qualitative account of stochastic acceleration.

The estimates given above are the upper limits on the emission frequency and typical energy of particles. They are reached only if the correlation length of the field fluctuations \( l_c \) is of the order of the radiative length

\[
l_r \sim c \frac{U}{dt} = \frac{r_e^2}{\gamma r_e}
\]  

(25)

which in the case of maximally efficient acceleration becomes \( l_r r_L \sqrt{\frac{\rho_L}{\rho_e}} \). If, on the other hand, the correlation length of the fluctuating electric field is much less than the radiative length, \( l_c \ll l_r \), the effective accelerating field will be lower by a factor \( \sqrt{L_c/L_r} \), which represents an effective RMS deficit in the coherence of the particle velocity and accelerating electric field.

Suppose that the acceleration is lower than the maximum acceleration (Eq. (18)) by a factor \( f < 1 \)

\[
\frac{dU}{dt} \sim f \omega_{eff} mc^2.
\]  

(26)

Then we would have a typical frequency of emission

\[
\omega_t \sim f c/r_e,
\]  

(27)

a typical energy

\[
< \gamma > \sim \sqrt{r_L r_e},
\]  

(28)

and radiative length

\[
l_r \sim \frac{r_L}{\sqrt{f}} \left( \frac{r_L}{r_e} \right)^{1/2}.
\]  

(29)

Suppose that the turbulent energy is concentrated near the smallest turbulent scale, which may be approximated by the relativistic skin depth \( r_s^{(r)} = \sqrt{< \gamma > c/\omega_p} \). (This will provide a lower limit to the efficiency of acceleration.) The ratio of the radiative length to the smallest scale of turbulence \( r_s^{(r)} \) then becomes

\[
\frac{l_r}{r_s^{(r)}} \sim \frac{1}{\gamma^{3/2}} \frac{r_e}{r_s r_e}.
\]  

(30)

The efficiency of the acceleration \( f \) can be estimated as \( f \sim \sqrt{\frac{l_s^{(r)}}{r_s}} \). Then (28) and (30) give

\[
\frac{l_r}{r_s^{(r)}} = \left( \frac{r_e r_L}{r_s^2} \right)^{2/7} = 4 \times 10^{-6}.
\]  

(31)

Since at the transition radius \( \frac{r_e}{r_s} = \frac{20}{e} \), the efficiency of the acceleration \( f \) becomes

\[
f = \left( \frac{r_e r_L}{r_s^2} \right)^{1/7} = \left( \frac{2e \Omega}{c} \right)^{1/7} = 2 \times 10^{-3}.
\]  

(32)

The typical energy of emitted photons in this case

\[
\epsilon \sim f h \frac{c}{r_e} = 80 \text{keV}
\]  

(33)

and the average energy \( < \gamma > \sim \sqrt{r_L r_e} = 5 \times 10^3 \).

Relation (33) gives the peak energy in the plasma frame. In the laboratory frame this should be multiplied by the Lorentz factor of the flow. Since the observed peak energy are around 500 keV, we infer that the bulk Lorentz factors are relatively small.
The low bulk Lorentz factors may reflect the fact that MHD acceleration usually is not very effective (Michel 1969).

We would like to stress that Eq. (22) was derived by assuming that the turbulent energy is concentrated on the skin depth scale. The presence of larger scale electric fields would increase the efficiency factor $f$, so Eq. (22) is a lower limit. The weak dependence of this lower limit on the magnetized rotator’s spin $\Omega$, its independence of the rotator’s magnetic field and thus independence of the local plasma density parameter, $\lambda$, are interesting properties of our model. Said another way, when this lower limit on $f$ applies, the energy of the emitted photons is independent of two out of three free parameters of the model ($B_0$ and $\lambda$), and only very weakly dependent on the third $\Omega^{1/7}$.

Next we confirm our earlier assumption that the emission resembles synchrotron emission not “jitter” emission. The typical nonrelativistic Larmor radius $r_L \sim 10^{-3} \text{cm}$ is much less than the smallest turbulent scale (given by the the relativistic skin depth $r_s^{(s)} = c\sqrt{\gamma_m}/\omega_p \sim 10^3 \text{cm}$, $< \gamma_m > = 5 \times 10^3$). Thus the particles will produce synchrotron-type radiation in a spatially varying magnetic field.

We can also verify that that the plasma in optically thin to Thompson scattering ($\tau_T = \sigma_T n_T l \sim 10^{-3}$) and to synchrotron selfabsorbtion at peak frequency (Pacholczyk Eq. 3.45).

### 4.4. Spectrum of turbulence

The actual spectrum of the relativistically strong EM plasma turbulence has not yet be determined, so we resort to general arguments. One possibility is a power-law Kolmogorov-type spectrum in with dominant turbulent energy on large scales, transferred to small scales. The problem is that relativistically strong EM turbulence may not satisfy the following necessary conditions required for by Kolmogorov-type turbulence: (i) the turbulent cascades should be local in $k$ space and (ii) there should be an inertial range in $k$ space where resistive damping is weak. Both these conditions may not be satisfied: strong turbulence may be nonlocal and damping of turbulent energy by plasma particles may be important.

Another possibility is a Phillips-type spectrum (Phillips 1958). Following Phyllips (1958) we can argue that at the point of overturn the profile of the wave becomes a discontinuous function (discontinuity of the zeroth order). The large wave number limit of the Fourier transform of such function will be $\propto k^{-1}$. Thus we can argue that the resulting spectrum of the EM turbulence will be $E_\parallel \propto k^{-2}$.

Assuming that $E_\parallel = Ck^{-2}$ and that the smallest scale of the turbulence is given by the relativistic skin depth, we find

$$E = \int d^3k E_\parallel \sim C \int dk = Ck_{max} = C \frac{1}{r_s}.$$  

In this case, the maximum power per range of $dk$ is $\propto k$. Most of the turbulent energy is indeed concentrated at the smallest turbulent scales $\sim r_s$. For Kolmogorov-type turbulence, $E_\parallel \propto k^{-11/3}$ so that most of the turbulent energy is concentrated at the large scales.

The turbulent plasma may also produce emission due to nonlinear conversion of the turbulent EM fields into escaping radiation. The typical frequencies will be of the order of the plasma frequency, $\omega_p \sim 10^9 \text{sec}^{-1}$, falling into the radio wave band. Another possibility is the production of radio waves by the overturning wave itself - an analogue of the supercontinuum generation in optical shocks (e.g. Zozulya et al. 1999).

### 4.5. Temporal evolution and spectra

The temporal evolution of the emission from the foaming region depends on the onset and temporal behavior of the electromagnetic turbulence. In the initial stages of the wave breaking the energy of the initially coherent EM wave may be deposited at large scales, comparable to the size of the breaking region ($\sim 10^7 \text{cm}$). After the initial energy injection, the energy cascades to smaller scales through turbulent interaction, reaching the smallest scale of the turbulence, the skin depth.

This turbulent cascade may or may not reach a quasistationary state. The typical time for the develop-

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3 Recall that in our model the bulk Lorentz factors do not require $\Gamma_{bulk} \geq 100$.

4 There are three free parameters: $\Omega$, $B_0$ and multiplicity $\lambda$.

5 Dettmann & Frankel (1996) argued that in relativistic hydrodynamic turbulence $\Gamma \propto r$, where $r$ is the scale param-

6 This is different from the breaking of surface waves on deep water, which form unstable cusps with a discontinuous derivative (discontinuity of the first order). In the case of surface waves the spectrum of the turbulence is $\propto k^{-4}$ (Phillips spectrum).
opment of the turbulent cascade maybe as large as the overturn time of the largest turbulent eddies, \( \sim 1 \) sec, if the turbulent cascade is local in \( k \) space, or as small as the typical microscopic interaction time, \( \sim 1/\omega_p = 10^{-8} \) sec, if the turbulent cascade is non-local in \( k \) space. The properties of the relativistically strong plasma turbulence have not been investigated yet, so we leave the question of the quasistationary state open and discuss observational features of both cases.

If it takes longer than the plasma synchrotron cooling time for the turbulent energy to cascade to the smallest scales, then the quasistationary turbulent spectrum may not be reached. In this case we expect that the observed emission from such non-stationary turbulence will be almost time independent until (and if) the turbulence cascades to the radiation length scale. The reason is that initially the coherence size of the electric fields is much larger than the radiation length so the initial particle acceleration is maximally efficient. Particles are accelerated to roughly the same energy (since the energy is radiation, not acceleration, limited), so that the emission spectrum is quasistationary and similar to the synchrotron spectrum from a mono-energetic particle distribution. When the turbulent cascade reaches the radiation length, the acceleration ceases to be maximally efficient. The particles are then accelerated in stochastic electric field, so that their spectrum will be power law with time dependent diffusion coefficient. The diffusion coefficient will be decreasing with time, acceleration will be less and less efficient and correspondingly the observed spectrum will be softening with time.

In the limit of very slow energy cascade, the radiation losses may drain the energy contained in the EM turbulence before the typical turbulent scale reaches the radiation length. In this case the spectrum of the particles and of the emitted radiation will resemble the spectra of synchrotron cooling sources.

If on the other hand the turbulent energy cascades to small scales faster than the plasma synchrotron cooling time, a quasi stationary distribution of electrons will be reached. Acceleration due to small scale electric fields will be stochastic while acceleration due to large scale electric fields will be of DC type. The particle acceleration will depended on the spectrum of the EM turbulence. If most of the energy is concentrated at large scales, as in Kolmogorov turbulence, then acceleration will be maximally efficient and the spectrum of particles will be approximately monoenergetic. If, on the other hand, most of the energy is concentrated at small scales, the acceleration will be stochastic.

The observed spectra of GRBs will depend on the distribution function of emitting electrons, which, in turn, depends on several unknown parameters like the prevailing type of acceleration (stochastic or DC-type), proper boundary and initial conditions (continuous ejection of particles at low or high energies or ”no flux” equilibrium states (Melrose 69, 71, 1980, Tademaru 71, Park and Petrosyan, Katz 1994)), relative importance of escape of highly energetic particles in the long duration bursts. The stationary ”no flux” equilibrium (Tademaru 71) will produce a relativistic Maxwellian distribution \( f(\gamma) \propto \gamma^\alpha e^{-\gamma/\gamma_0} \), with \( \alpha > 2 \), while ”reflecting boundaries” condition of Melrose (69, 71, 1980) will produce power law at \( \gamma > \gamma_0 \). DC-type acceleration will tend to produce strongly peaked (monoenergetic) distribution, while stochastic acceleration will lead to spectral broadening. In addition, if particle escape is important, the resulting spectra will be powerlaw-like, depending on the energy dependence of the escape probability (Melrose 1980). This variety of possible spectra may serve well to explain the unusual variety of GRB’s spectra, which are often powerlaws at small and large energies and sharply peaked near the spectral break (Band et al. 1993). We leave a more detailed consideration of spectral properties for a subsequent work.

5. Regeneration of magnetic fields and afterglows

The interface between the accelerated plasma and the ambient medium should lead to the formation of an external shock at large distances from the magnetized rotator. This external shock could correspond to that inferred to be responsible for GRB afterglows (c.f. Meszaros & Rees 1993; Sari et al. 1998; Piran 1999; Wijers & Galama 1998, Frail et al.(1999), Gruzinov et al.(1999)), at large distances. To fit the afterglows with the synchrotron emission of relativistic particles large magnetic fields, of the order of equipartition are required. Generation of such large magnetic fields have presented an unsolved problem (Gruzinov(1999)) in the blast wave model.

7 In this respect we also note that the ”plasma turbulent reactor” (refs**), which bears some similarity to our model problem, produces a power-law spectrum \( p^{-3} \), consistent with the observed spectral index \(-2\).
In the PFDO model there is a possibility of regeneration of magnetic fields due to interaction of LAEMW with plasma. It has been known that interaction of LAEMW in the form of lasers with plasma produces both large and small scale magnetic field (e.g., Stamper et al. 1975; 1978, Sudan 1993, Max 1980). Most commonly, magnetic fields are generated on skin depth scales (Pegoraro et al. 1997). The typical regenerated magnetic fields may be as large as initial magnetic field $B \sim 10^6$ G. It has been argued that such small scale quasistationary magnetic fields may explain the polarization observed in the afterglows (Medvedev 2000). Alternatively, the magnetic fields may be also subject to coalescence instability (Finn & Kaw 1977, Sakai et al. 199), in which randomly oriented magnetic islands merge forming large scale structures.

There is also a possibility of generating large scale magnetic fields, with correlation length of the order of the size of inhomogeneities of the “of the laser beam” by inverse Faraday effect (Berezhiani et al. 1997), which may generate fields of the order

$$B \sim \frac{\gamma_0 m c^2}{e R_{1\perp}}$$  \tag{35}$$

where $\gamma_0 \sim \nu_0$ is the typical Lorentz factor of particles accelerated by a EM pulse with a transverse dimension $R_{1\perp}$, in our case the transverse size of the foaming region.

We mention these possibilities here as a flag for future work, and to highlight the need to further investigate possible analogies between field amplification in GRB and that of laser driven plasmas. The connection is that both plasma systems may be driven by LAEMW.

6. Conclusion

The two main points of this work may be summarized as follow: (i) Poynting flux dominated outflows, which so far has been invoked only in the context of central engines, possesses internal instabilities which may also explain radiation generation, (ii) emission from stochastically accelerated particles in turbulent electromagnetic fields, when electric fields are as important as magnetic fields in the acceleration and radiation process, is a viable mechanism for the GRB emission.

Our approach may also naturally explain the result that the peak of GRB emission varies only over a small range of value from burst to burst. (Brainerd 1994, 1997). For our maximally efficient mode of acceleration, the peak, when boosted into observer frame, becomes $\epsilon_{\text{max}} \sim f h c / r_e \Gamma_{\text{bulk}}$. If this maximally efficient regime were operating, the bulk gamma-factor $\Gamma_{\text{bulk}}$ would not be much larger than 10. The location of this peak is a very weak function of the parameters of the underlying rotator, $f \sim \Omega^{1/7}$ and is independent of the progenitor’s magnetic field.

This can be compared to phenomenological internal shock models in which the peak energy is proportional to $\gamma_{\text{min}}^2 \omega_p \Gamma_{\text{bulk}}$, where $\gamma_{\text{min}}$ is the low energy cutoff to the electron power law spectrum. All three quantities here are usually taken to be free parameters. It seems unlikely for their combination to be almost constant from burst to burst, and within each pulse of a given burst unless some physics dictates this to be the case. We have tried to add some of this physics. In our approach the remaining free parameter is the Lorentz factor of the bulk flow.

PFDOs may also resolve the problem of the magnetic field generation since in the Poynting flux driven outflows the large magnetic fields are supplied by the source. These fields do not necessarily have to be generated in the external shock for the afterglow since the shock is a current sheet through which outflow and ambient particles mix (Smolsky & Usov 2000). At the same time, however there exists an unexplored analogy between field generation mechanisms in laser driven plasmas and GRB outflows that will have to be understood to determine the scale and structure of the outflow magnetic field.

Other observational properties of the GRBs that may be explained in our framework: (i) GRBs show no correlation between the spectra and other micropulse characteristics (e.g., Lee & Bloom 2000) - this is a direct consequence of the turbulent EM acceleration which produces photons with the frequency $\sim c/r_\perp$; (ii) the fact that pulses peak earlier at high frequencies and that bursts have shorter duration at higher energies may be due to the initial development of the turbulence: if the turbulence develops from large scales to small scales then initially the acceleration may be more effective since it is due to large scale electric fields (with coherence length larger than the radiative length). In this case initially particles are accelerated to the limiting energies $\gamma \sim \sqrt{r_L / r_w}$ while later, when the coherence length of the turbulence becomes smaller than the radiative length, the particle distributions soften and radiation spectra emanates...
at lower frequencies. The typical time for such a cascade should be of the order of the "vortex overturn time" on the largest scale of the turbulence, which may be on the order of seconds. (iii) The absence of correlation between the pulses and overall burst characteristics, interpreted as arising from random and independent emission episodes (Lee & Bloom 2000), is natural in our model since each wave overturn happens independently. (iv) The temporal characteristics of the microbursts, FREDs, and the hard-to-soft spectral evolution is a consequence of synchrotron cooling of the reservoir of energy released during wave overturn. (v) Composite structure of GRBs (a burst being a sum of many independent emission events) (Stern & Svennson 1996) naturally follows from the model - each overturning region produces an independent microburst. (vi) The average power density spectrum of GRBs is well described as being due to selfsimilar turbulent-type process near marginal stability (Stern 1999). This is reminiscent of the cellular automata model of solar reconnecting regions (Lu & Hamilton 1991) and may be related to the reconnecting "foaming" regions in our model.

We also would like to point out that the particular mechanism of the wave instability, the wave overturn during MHD-wave transition, may not necessarily be the only one. Other plasma instabilities may contribute to the generation of EM turbulence.

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Fig. 1.— Profiles of (a) parallel momenta and (b) density in a wave. The profile of the wave is 
\[ a = v_0 \arctan(x - v_g t) \] 
with \( v_g = 1/\sqrt{1 + \mu n/\gamma} \). Density parameter \( \mu = 0.1 \). Thin lines correspond to subcritical amplitude \( v_0 = 2.4 \), thick line correspond to overcritical amplitude \( v_0 = 2.5 \).