Laser-guided focusing of intense relativistic electron beams for fast ignition

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(Received 25 November 2003; accepted 10 May 2004; published online 12 July 2004)

As an alternative to the fast ignition by petawatt lasers of small spherical deuterium–tritium (DT) targets compressed to thousand times solid density, the fast ignition by terawatt relativistic electron beams of thin cylindrical (or conical) DT targets, compressed to several ten times solid density and magnetized to 10^8 G through a high current discharge along the cylindrical axis of the targets, has been proposed. One problem of this approach is the guiding and focusing of the relativistic electron beam onto the target. It is proposed to transport the beam through a laser-triggered ionized channel in a low density background gas, or by letting it propagate along the surface of a thin wire crossing the diode gap, with the final focusing done by repulsive image currents in a conducting convergent cone. A second problem is the stopping of the electron beam in the target. This can hopefully be done by a combination of classical electron stopping power, the electrostatic two-stream instability and collisionless shocks in the presence of a strong perpendicular magnetic field. © *2004 American Institute of Physics.* [DOI: 10.1063/1.1768178]

I. INTRODUCTION

Fast ignition by petawatt lasers[1] of small spherical deuterium–tritium (DT) targets requires 1000-fold solid density compression of these targets. For the compression to these ultrahigh densities, pulsed lasers,[2,3] light[4] and heavy[5,6] ion beams, but also soft x rays from electric pulse power imploded wire arrays,[7] have been proposed.

As a potentially much less expensive alternative the fast ignition by terawatt relativistic electron beams of cylindrical (or conical) DT targets, compressed to several ten times solid density and magnetized in excess of 10^8 G through a high current discharge along the axis of these targets, has been proposed.[8,9] There two nested magnetically insulated transmission lines are used; the outer line carrying a high current–lower voltage pulse for compression and α-particle confinement, and the inner line carrying a high voltage–lower current pulse for fast ignition (see Fig. 1).

The relativistic electron beam for fast ignition is there emitted from the cathode at the end of the inner transmission line on the left of Fig. 1, and projected on the conical DT target to the right. The required $\rho r > 1$ g/cm^2 for the ignition of spherical DT assemblies is here replaced by $\rho z > (1/3)$ g/cm^2, valid for a one-dimensional burn wave, where z is the minimum length of cylinder or height of the cone.

A typical example for the outer line is $I_0 = 10^7$ A, $V_0 = 10^6$ V lasting 10^{-8} s, and for the inner line, $I = 3 \times 10^5$ A, $V = 10^7$ V lasting 10^{-9} s with a total energy of about 100 kJ.

In the detail of Fig. 1 shown in Fig. 1(a), the current of the outer high current line is wound in a spiral around the return current conductor of the inner high voltage line enhancing the magnetic insulation of the inner line.

There are three advantages for this alternative fast ignition scheme.

(1) Electric pulse power is by orders of magnitude less expensive than pulse power by laser beams or particle beams.

(2) With a strong magnetic field permitting smaller DT densities, the power required for fast ignition is reduced in direct proportion of the DT target density from 10^{15} W down to a few 10^{12} W, the latter attainable with electric pulse power.

(3) Because of plasma confinement by the strong magnetic field, the energy required for fast ignition is reduced from ≥ 100 kJ to ≥ 1 kJ, and the energy for compression from several megaJoules to ≥ 100 kJ.

It is worthwhile to mention that this concept evolved from an older idea where a shear flow stabilized dense z pinch below the DT ignition temperature is ignited at one end by a pulsed laser beam, with a thermonuclear detonation wave propagating with supersonic speed along the pinch discharge channel.[10]

II. FOCUSING AND LASER GUIDANCE OF THE INTENSE RELATIVISTIC ELECTRON BEAM

The feasibility of the proposed alternative concept depends on the ability to transport and focus the relativistic electron beam emitted from the field emission cathode onto the small cross section at the end of the DT target cylinder. Naively one may think that by the Child–Langmuir law the buildup of electric space charge in between the field emission cathode and DT target anode prevents a large electron beam flow inside the diode. However, at the very high field strengths and in the presence of a background gas this law is invalid, because there the current is determined by electric breakdown. With a pulsed laser beam of modest energy making an ionized trail in the background gas, one can guide the electron beam through a thin channel in the background gas (see Fig. 2). For complete space charge neutralization of the

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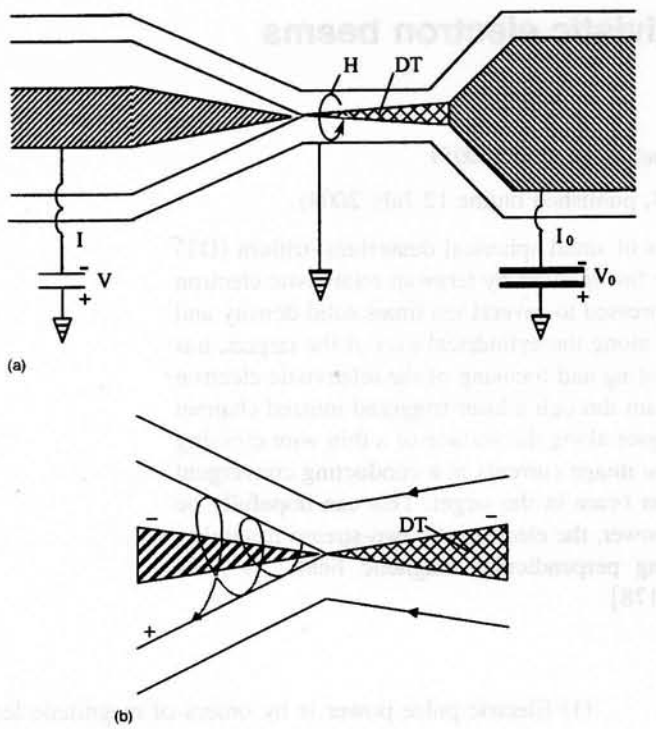

FIG. 1. (a) The two nested magnetically insulated transmission lines with the DT target ignited in the focus of both lines. In our example we have chosen for the inner line a voltage of $V=10^7$ V and current $I=3\times10^5$ A; and for the outer line $V_0=10^6$ V and $I_0=10^7$ A. (b) The two nested transmission lines with improved magnetic insulation of the inner high voltage line by a helical winding of the high current flow of the outer line over the return current conductor of the inner line.

electron beam the number density of the background gas must be equal to the number density of the electron beam. With the electron beam propagating through the ionized channel, and with a large axial electric field inside the diode, the conditions for electron run-away can be established, with the result that the electron beam impacts the anode with its full kinetic energy. Instead of letting the beam propagate through an ionized channel in a low density background gas, one can also propagate it along a thin wire crossing the diode gap. Final beam focusing onto a precisely prescribed spot on the anode for ignition can be achieved by projecting the

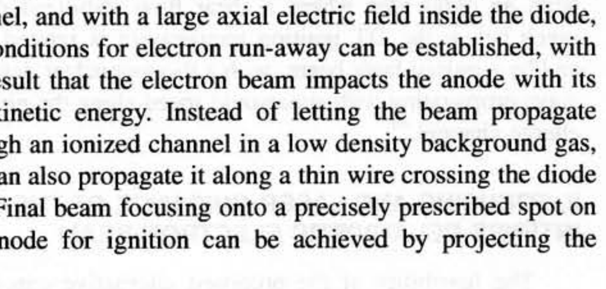

FIG. 2. Laser guided run away relativistic electron beam in space charge neutralizing background gas with laser beam initiating breakdown. ℓ guiding laser beam, C cathode, A anode, RC return current conductor, e relativistic electrons.

beam into a convergent conducting cone attached to the DT target, as it is realized in the diode configuration of Figs. 1 and 2.

In addition to electron run-away three other effects have to be considered.

(1) The thermomagnetic Nernst effect in the beam channel.

(2) Beam heating by Coulomb collisions with the background gas.

(3) Beam cooling by synchrotron radiation.

As will be shown, in a singly ionized background plasma the thermomagnetic Nernst effect establishes a magnetic field of the same strength as the beam would have in vacuum, compensating the return current in the background plasma. Because of the strong azimuthal magnetic field set up in the electron beam, the electrons oscillate in a direction perpendicular to the beam losing energy by synchrotron radiation, and for a reasonably large γ values ($\gamma=1/\sqrt{1-v^2/c^2}$), at a rate larger than the heating rate by Coulomb collisions. For very large γ values, this effect can lead to a substantial lowering of the beam emittance, of importance, if the fast ignition is done with an electron beam from a high energy particle accelerator.

III. LIMITS ON BEAM FOCUSING

The smallest possible beam radius is determined by Liouville's theorem (m electron rest mass)

$$mrv=\gamma m r_0 c, \tag{1}$$

where v is the nonrelativistic velocity of the electrons in the field emission cathode, determined by the electron Fermi energy $E_F=(1/2)mv^2=10$ eV, and r the initial beam radius. Furthermore, $\gamma\simeq E/mc^2$, where E is the energy of the relativistic electrons. One thus obtains from Eq. (1) for the smallest beam radius r_0

$$\frac{r_0}{r}=\sqrt{\frac{2}{\gamma}}\sqrt{\frac{E_F}{E}}. \tag{2}$$

If $E=10$ MeV with $\gamma\simeq20$, one finds that $r_0/r\simeq3\times10^{-4}$, which means that focusing down to $r_0\simeq10^{-2}$ cm is possible even if r is quite large.

IV. ENERGY REQUIREMENTS FOR THE LASER TRIGGERED DISCHARGE CHANNEL

The formation of a laser triggered channel for a pinch discharge in a dense gas or solid matter was proposed by Tidman,[11] but it is also used to ignite sparks in laser triggered gas switches.[12] We propose here to use the same technique to make an electron run away discharge channel in the gas filled high voltage diode.

How this can be done is shown in Fig. 2, where the laser beam passes through a small hole in the center of the inner high voltage cathode. To avoid diffraction, the wave length of the laser light must be small compared to the diameter of the hole in the high voltage cathode shown in Fig. 2, a condition well satisfied for a cross section of the hole of the order 10^{-2} cm^2.

In order to ionize the channel, it must be heated to a temperature of ~ 10 eV$\sim 10^5$ K. For a number density of the background gas equal to $n \simeq 10^{17}$ cm^{-3} and an optical cross section for absorption of the laser light which (as for dark glass) is assumed about equal to $\sigma_{\rm opt} \sim 10^{-20}$ cm^2, the optical path length is $\approx 10^3$ cm large enough to uniformly heat a channel a few centimeters long. With a cross section of 10^{-3} cm^2 and a length of 1 cm the channel contains $N = 10^{17} \times 10^{-3} = 10^{14}$ atoms to be heated to 10^5 K requiring an energy of $3NkT \sim 4 \times 10^3$ erg$= 4 \times 10^{-4}$ J. With an optical path length of $\sim 10^3$ cm, in the channel only, 0.1% of the laser energy is absorbed. The laser energy must be therefore 10^3 times larger or about 0.4 J. If delivered in 10^{-9} s the laser power would be 400 MW.

V. BEAM FORMATION BY ELECTRON RUN-AWAY

The condition for electron run-away in a plasma is[13]

$$eE\lambda \gg (1/2)mv^2 = (3/2)kT, \tag{3}$$

where e and m are the electric charge and mass of an electron, E the applied electric field and

$$\lambda = 1/n\sigma_c, \tag{4}$$

the collision mean free path in the background gas of number density n with the collision cross section

$$\sigma_c = \frac{4\pi}{9}\frac{Z^2 e^4}{(kT)^2}. \tag{5}$$

Inserting Eq. (4) into Eq. (3) and using Eq. (5) one has

$$2.2 \times 10^9 \frac{TE(\text{V/cm})}{Z^2 n} \gg 1. \tag{6}$$

In our example $n \simeq 10^{17}$ cm^{-3}, and we assume $Z=1$ (singly ionized background gas). With a distance between the cathode and anode about ~ 1 cm, and a voltage of $\sim 10^7$ V, one has $E \sim 10^7$ V/cm. For these numbers the inequality (6) is well satisfied for temperatures larger than 10^3 K, and a 10 MeV relativistic electron beam can be easily established. For the axial electric field to be able to accelerate the electrons the field penetration depth must be larger than the beam radius at its smallest diameter, which is of the order 10^{-2} cm. The penetration depth is given by

$$\delta = c/\omega_p, \tag{7}$$

where $\omega_p = \sqrt{4\pi n e^2/\gamma^3 m}$ is the longitudinal beam plasma frequency with $\gamma^3 m$ the longitudinal electron mass. For $n = 10^{17}$ cm^{-3} and $\gamma = 20$ one finds that $\delta \simeq 0.5$ cm which is large enough for the externally applied electric field to penetrate into the entire electron beam plasma column. For an oblique direction of the electric field relative to the electron beam, one has to take the transverse electron mass γm instead, whereby the penetration depth would be at most γ times smaller, for $\gamma = 20$ about equal to 2.5×10^{-2} cm, still large enough.

VI. THE THERMOMAGNETIC NERNST EFFECT FOR THE GAS-EMBEDDED ELECTRON BEAM

The importance of the Nernst effect in rapidly flowing wall confined plasma configurations[14,15] and for the focusing of plasma jets[16] has been previously studied and is similar to the situation of a gas-embedded electron beam. As there, the temperature gradient between the hot electron-beam-carrying plasma channel and the surrounding cooler gas leads to a Nernst current with the density[13]

$$\mathbf{j}_N = \frac{3knc}{2H^2}\mathbf{H} \times \text{grad}\, T, \tag{8}$$

where $\mathbf{H}$ is the magnetic field, with $\mathbf{j}_N$ connected to $\mathbf{H}$ by Maxwell's equation

$$\frac{4\pi}{c}\mathbf{j}_N = \text{curl}\, \mathbf{H}. \tag{9}$$

In a cylindrical coordinate system with z directed along the beam current one has

$$\mathbf{j}_N = -\frac{3knc}{2H}\frac{dT}{dr} \tag{8$'$}$$

and

$$\frac{4\pi}{c}\mathbf{j}_N = \frac{1}{r}\frac{d}{dr}(rH), \tag{9$'$}$$

where $H = H_\varphi$. Elimination of $\mathbf{j}_N$ from (8′) and (9′) leads to

$$-6\pi kn\frac{dT}{dr} = \frac{H}{r}\frac{d}{dr}(rH). \tag{10}$$

The magnetic force density of the Nernst current is

$$\mathbf{f} = \frac{1}{c}\mathbf{j}_N \times \mathbf{H} = \frac{3}{2}\frac{nk}{H^2}(\mathbf{H} \times \nabla T) \times \mathbf{H}, \tag{11}$$

or since ∇T is perpendicular to $\mathbf{H}$,

$$\mathbf{f} = \tfrac{3}{2} nk \nabla T. \tag{12}$$

With $p = 2nkT$ for a singly ionized plasma and Eq. (11) the magnetohydrodynamic equilibrium equation

$$\nabla p = \mathbf{f} \tag{13}$$

becomes

$$2nk\nabla T + 2kT\nabla n = \tfrac{3}{2} nk \nabla T \tag{14}$$

or

$$\frac{\nabla n}{n} = -\frac{1}{4}\frac{\nabla T}{T} \tag{15}$$

hence

$$Tn^4 = \text{const}. \tag{16}$$

We then set

$$n = n_0 T_0^{1/4}/T^{1/4}, \tag{17}$$

where T_0 is the maximum temperature at the center of the beam plasma channel where $n = n_0$. For constant current distribution in the channel one has

$$H = H_0 \frac{r}{r_0}, \tag{18}$$

where H_0 is the field strength at the channel boundary with radius r_0. Inserting Eqs. (17) and (18) into Eq. (10) one obtains

$$-6\pi k n_0 T_0^{1/4} \frac{dT}{T^{1/4}} = \frac{H_0}{r_0} d(rH), \tag{19}$$

which by integration from T_0 to T on the left side, and from $Hr=0$ to $Hr=H_0 r_0$ on the right side, yields

$$\frac{H_0^2}{8\pi} = n_0 k T_0, \tag{20}$$

which is the magnetohydrostatic equilibrium for constant current distribution. It thus follows that the magnetic field set up around the plasma-embedded electron beam is the same as it would be for the beam in vacuum.

For $n_0 = 10^{17}\ \mathrm{cm}^{-3}$, $I = 3\times 10^5$ A and $r_0 = 10^{-2}$ cm, hence $H_0 = 6\times 10^6$ G, one obtains from Eq. (17) $kT_0 = 1.4 \times 10^{-5}\ \mathrm{erg} \simeq 9\times 10^6\ \mathrm{eV} \simeq 10$ MeV, consistent with a 10 MeV electron beam. With $n_0 r_0^2 = nr^2$ and $H_0^2 r_0^2 = H^2 r^2$, this result remains unchanged for a beam radius $r > r_0$.

VII. BEAM HEATING AND COOLING

During its transport through a background gas from the field emitting cathode to the anode, there is transverse electron beam heating by the Coulomb collisions with the background gas, but because of the large transverse electron velocity, the beam is also cooled by the emission of synchrotron radiation. Beam heating by Coulomb collisions is given by[17]

$$\frac{dE}{dt} = 4\pi n \frac{e^4}{mc} \log \Lambda, \tag{21}$$

where $\Lambda = r_0/r_e$, r_0 the beam radius, and $r_e = e^2/mc^2$ classical electron radius. dE/dt is the energy loss rate for one electron. With the same number density n in the beam and the background hydrogen plasma, and with the beam current given by $I = nec\pi r_0^2$ one obtains

$$\frac{dE}{dt} = \frac{ce^2}{r_0^2} \frac{I}{I_A} \log \Lambda, \quad I_A = mc^3/e = 17\,000\ \mathrm{A}. \tag{22}$$

Beam cooling is caused by the synchrotron radiation of the electrons radially oscillating in the magnetic potential well of the beam. These synchrotron radiation losses are[18]

$$P_e = \frac{2}{3} \frac{e^2 \dot{v}_\perp^2}{c^3} \gamma^4. \tag{23}$$

The magnetic field inside the beam is

$$H = \frac{2I}{r_0 c} r \tag{24}$$

and the radial restoring force on a beam electron ($\beta = v/c$)

$$F = -e\beta H = -\frac{2e\beta I}{cr_0^2} r = -\frac{2eI}{cr_0^2} r. \tag{25}$$

With the transverse electron mass equal to γm one has for the radial electron oscillations

$$\gamma m \frac{d^2 r}{dt^2} - F = 0 \tag{26}$$

or

$$\frac{d^2 r}{dt^2} + \omega^2 r = 0, \tag{27}$$

where $\omega^2 = 2eI/\gamma m c r_0^2 = (2/\gamma)(c/r_0)^2 (I/I_A)$. For a harmonic oscillator one has $\dot{v}_\perp = \langle (d^2r/dt^2) \rangle = -(1/2)\omega^2 r_0$ and hence

$$\mathbf{P}_e = \frac{2}{3} \frac{e^2 c}{r_0^2} \left(\frac{I}{I_A}\right)^2 \gamma^2. \tag{28}$$

Since $\omega^2 = (2/\gamma)(c/r_0)^2 I/I_A = 4\pi n e^2/\gamma m$, the electron beam plasma remains transparent.

The beam emittance gets smaller if

$$\mathbf{P}_e > dE/dt. \tag{29}$$

With Eqs. (22) and (29) this happens if

$$\gamma^2 (I/I_A) > (3/2) \log \Lambda. \tag{30}$$

The condition for the beam not to pinch itself off by its own magnetic field requires that

$$H^2/8\pi < \gamma n m c^2, \tag{31}$$

which means that

$$I < \gamma I_A. \tag{32}$$

For the given example $\gamma = 20$, $I \sim 3\times 10^5$ A, $I \simeq \gamma I_A$ one obtains from Eq. (30)

$$\gamma^3 > 3 \log \Lambda. \tag{33}$$

For $r_0 \simeq 10^{-2}$ cm, $r_e \simeq 10^{-13}$ cm, $r_0/r_e \simeq 10^{11} = \Lambda$, and $\log \Lambda \simeq 25$, one finds $\gamma > 5$. This means that for $\gamma = 20$, beam heating is overcome by beam cooling. Furthermore, with the beam collapse time given by

$$\tau \simeq \frac{r_0}{c\gamma^3} \left(\frac{r_0}{r_e}\right)^2, \tag{34}$$

one finds with $r_0 \sim 10^{-13}$ cm, $r_0/r_e \simeq 10^{11}$, $\gamma = 20$ that $\tau \sim 10^{-5}$ s, much too long to be of any significance. It thus follows that except for large γ values, beam heating (and cooling) can be ignored.

VIII. FINAL BEAM FOCUSING

Final beam focusing down to a beam radius $r_0 \simeq 10^{-2}$ cm, as required for ignition, can be done by projecting the beam into a convergent cone with a conducting wall. There by image currents in the wall the beam can be deflected as shown in a nice photograph published by Physics International,[19] reproduced here as Fig. 3.

With the electron beam stagnation pressure

$$p = \gamma n m c^2 \tag{35}$$

and the magnetic beam pressure

$$p_H = H^2/8\pi, \tag{36}$$

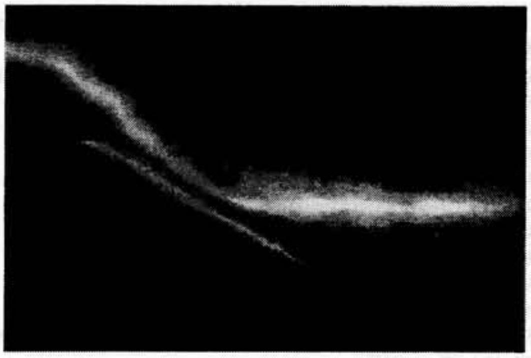

FIG. 3. Deflection of intense relativistic electron beam by image currents in a conducting wall, with the beam going from the right to the left (photograph by Titan Pulse Sciences Division, formerly, Physics International, San Leandro, California).

where $H=2I/rc$ and $I=\pi r^2 nec$, one obtains for the β value of the beam

$$\beta=p/p_H=2\gamma I_A/I. \tag{37}$$

Then, if α is the angle of incidence between the beam and the conducting wall, beam deflection requires that

$$p\sin\alpha<p_H \tag{38}$$

or that

$$\alpha<\arcsin(1/\beta). \tag{39}$$

For $I\simeq\gamma I_A$, hence $\beta=2$, one finds that $\alpha<30°$, requiring a slender cone.

IX. TARGET HEATING

For target heating three processes must be considered.

(1) Heating by classical electron beam stopping in compressed DT.

(2) Heating by the electrostatic two-stream instability.[20–22]

(3) Heating by the formation of a collisionless shock in the presence of a strong transverse magnetic field.[23,24]

With the DT encapsulated in a thin metallic liner serving the conductor of the outer high current transmission line, the initial DT temperature should be comparable to the liner temperature, which as for exploding wires is estimated to be a few 10^6 K. With a current $I_0\simeq10^7$ A flowing over a cylindrical liner with a radius equal to $r_1=10^{-2}$ cm, the magnetic field at the liner surface is $H=2\times10^8$ G with a magnetic pressure equal to $H^2/8\pi\simeq1.6\times10^{15}$ dyn/cm^2. Equating this pressure with the DT plasma pressure inside the liner $p=2nkT$ and $T=3\times10^6$ K, one finds that $n\simeq2\times10^{24}$ cm$^{-3}=40n_0$, where $n_0=5\times10^{22}$ cm^{-3} is the number density for solid DT or $\rho\simeq8.4$ g/cm^3.

The classical electron stopping power range of relativistic electrons with energy E expressed in MeV is approximately given by

$$\lambda=\frac{1}{\rho}(0.543E-0.16)\ \text{(cm)}, \tag{40}$$

for 10 MeV electrons and $\rho=8.4$ g/cm^3, one has $\lambda\simeq0.6$ cm. According to the virial theorem applied to the beam electrons in the radial direction one has

$$(1/2)v_\perp^2/c^2=p_H/p=1/\beta, \tag{41}$$

with $\beta\simeq1$, $v_\perp\simeq c$. The electrons there move under an angle of 45° inside the beam, repeatedly reflected by the beam boundary. This increases their path length inside the beam by the factor $\sqrt{2}$, and their stopping range along the beam axis by the factor $1/\sqrt{2}$, from 0.6 cm down to 0.4 cm. This is still short enough to satisfy the $\rho z>(1/3)$ g/cm^2 condition for the ignition of a one-dimensional burn wave by a relativistic electron beam.

Much shorter stopping distances seem to be possible with the electrostatic two-stream instability, but the growth rate of this instability is reduced by the large translational motion of the electrons, having the same effect a hot electron beam would have. For a cold beam the range is

$$\lambda=\frac{1.4c\gamma}{\omega_p\varepsilon^{1/3}},\quad \varepsilon=\frac{n_b}{n}, \tag{42}$$

where $n_b\simeq10^{17}$ cm^{-3} is the number density of the electron beam and $n\simeq2\times10^{24}$ cm^{-3} the number density of the 40-fold compressed solid DT target, with $\omega_p\simeq8\times10^{16}$ s^{-1} the electron plasma frequency of the target. For 10 MeV electrons $\gamma\simeq20$, and one obtains $\lambda\simeq10^{-3}$ cm.

The thickness of a collisionless shock is of the order of the ion gyroradius, which for a DT plasma at the temperature $T\sim10^8$ K and a magnetic field $H\sim10^8$ G is of the order 10^{-4} cm.

All three effects, classical electron beam stopping power, electrostatic two stream instability, and collisionless shock formation support the idea of fast ignition by a relativistic electron beam in the described configuration.

X. CONCLUSION

The enormous cost of the laser fusion (with direct or indirect drive) makes it prudent to look out for other potentially less expensive inertial confinement fusion schemes.

One conceptually very simple scheme is impact fusion.[25,26] In comparison to laser fusion it is even better in solving the so called "stand off" problem of thermonuclear microexplosions. And in comparison to the poor laser efficiency, it has the high efficiency of an electric motor with the acceleration of the macroparticles done by a magnetic traveling wave accelerator. Furthermore, the robustness of such a macroparticle accelerator promises a long life time, important for the economy of a fusion power plant. But with all its advantages it is not cheap, requiring a linear accelerator more than 10 km long or a circular accelerator with a comparable circumference.

Replacing lasers with heavy ion beams leads to some cost reduction and a better "rep rate" capability. For indirect drive schemes a tenfold cost reduction may be possible by replacing lasers (or heavy ion beams), with electric pulse power imploded arrays of thin wires, producing a burst of soft x-rays for target compression, but with the yield of the thermonuclear microexplosion in the gigaJoule range, the stand off problem is here quite serious, destroying the transmission line after each microexplosion, with no simple solution in sight.

The concept of a double high current–lower voltage magnetically insulated transmission line for confinement and a high voltage–lower current magnetically insulated transmission line for fast ignition, is the only concept known to this author with the potential of a high gain–low yield thermonuclear microexplosion, permitting a comparatively small yield, greatly easing the stand off problem. It is also the only known concept with the potential for an up to 1000-fold cost reduction in comparison to laser fusion or other with laser fusion competing inertial confinement schemes. The most important problem here is the guiding focusing of the relativistic electron beam onto the thin cylindrical DT target, with the stopping of the electron beam coming in second. But with the power for fast ignition of a lower density–highly magnetized DT target reduced 1000-fold from petawatts to terawatts, a few kiloJoule terawatt laser could here be used instead of a relativistic electron beam, if the fast ignition with a relativistic electron beam should turn out to be too difficult.

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