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A primer for electroweak induced low-energy nuclear reactions

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Abstract. Under special circumstances, electromagnetic and weak interactions can induce low-energy nuclear reactions to occur with observable rates for a variety of processes. A common element in all these applications is that the electromagnetic energy stored in many relatively slow-moving electrons can – under appropriate circumstances – be collectively transferred into fewer, much faster electrons with energies sufficient for the latter to combine with protons (or deuterons, if present) to produce neutrons via weak interactions. The produced neutrons can then initiate low-energy nuclear reactions through further nuclear transmutations. The aim of this paper is to extend and enlarge upon various examples analysed previously, present order of magnitude estimates for each and to illuminate a common unifying theme amongst all of them.

Keywords. Nuclear transmutations; low-energy nuclear reaction; electroweak.

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1. Introduction

In [1–6], diverse physical processes were considered and detailed computations on them performed to reveal conditions under which, in each case, the end result is a low-energy nuclear reaction (LENR) induced by electroweak interactions. Even though weak interactions are an integral part of the Standard Model of fundamental interactions, which unifies the electromagnetic force with the weak and the strong (nuclear) forces, nonetheless low-energy applications of weak interactions in condensed matter devices are novel and hence unfamiliar. All the existing condensed matter devices are essentially of electromagnetic origin. There are sound reasons for the latter circumstance. Charged particles in condensed matter (electrons or ions) normally possess low kinetic energies (typically a few eV or less) and yet they can trigger substantial electromagnetic processes which one can usefully harness. In sharp contrast, for an electron to undergo a weak interaction, say with a proton in an ion and produce a neutron, MeV range of energies is needed (because the neutron is heavier than the proton by about 1.3 MeV and hence there is an energy 'threshold' which must be overcome). It follows then that for neutron production (and subsequent nuclear transmutations) via weak interactions, special conditions in condensed matter systems must be found which accelerate an electron to MeV range of energies. Successful avenues to accomplish precisely this purpose were described in the papers quoted above and some others are explored here. The present paper is devoted to delineating the unifying features and to an overall synthesis of these rather different collective processes. We discuss the phenomenon of long-range weak isospin oscillations in a material due to electromagnetism. We also estimate the total flux of positrons and protons from a solar flare. Through the latter, we compare our flux of high-energy muons detected in a solar flare at CERN.

In [1], metallic hydride surfaces on which plasma oscillations exist were analysed. It was shown that the collective plasma oscillations on the surface could contribute some of their electric energy to an electron so that the following reaction was kinematically allowed:

$$W_{\rm electric} + e^- + p \rightarrow \; n + \nu_{\rm e}. \tag{1}$$

The relevant scale of the electric field $\mathcal{E}$ and the plasma frequency Ω needed to accelerate electrons to trigger neutron production is found to be

$$\frac{c\mathcal{E}}{\Omega} = \left(\frac{mc^2}{e}\right) \approx 0.5 \times 10^6 \ {\rm V}, \tag{2}$$

where c is the speed of light, m is the mass and $(-e)$ is the charge of the electron.

The particular condensed matter environment leads in this case to ultracold (that is, ultralow momentum) neutrons. These ultracold neutrons produced, have extraordinarily large nuclear absorption cross-sections and thus a high probability of producing nuclear transmutations and an extremely low probability of neutrons escaping beyond micron scale and smaller surface regions in which they are formed. There is also a high suppression in the production of high-energy γ-rays [2]. For such metallic chemical cells, comprehensive calculations of the rates of LENRs [3] were made which confirmed a robust production of new elements.

In [4], a magnetic analog of the above (the so-called exploding wire problem) was analysed. We found that a strong electric current carrying wire [7] could – under suitable conditions – channel the collective magnetic energy sufficiently once again to excite a certain fraction of electrons to undergo the weak interaction process

$$W_{\rm magnetic} + e^- + p \rightarrow \; n + \nu_{\rm e}. \tag{3}$$

The scale of current required here was shown to be of the order of the Alfvén current I_0

$$I_0 = \left(\frac{mc^2}{e}\right)\left(\frac{4\pi}{R_{\rm vac}}\right) \approx 17 \ {\rm kA}. \tag{4}$$

Observation of copious neutrons in exploding wire experiments is by now legion [8–13]. Experimentally, neutron production has also been confirmed for lightning [14], the big exploding wire in the sky, where typical currents are about 30 kA and higher.

Quite recently [5], another application of the magnetic mode inducing LENR has been made to unravel the mystery surrounding the observed particle production and nuclear transmutations in the solar corona and the solar flares [15–23]. Spectacular pictures of flux tubes are now available [24] showing giant magnetic flux tubes exiting out of one sunspot and entering into another. We showed theoretically [5] how these could lead to steady LENR. In fierce solar flares, we found that as a flux tube disintegrate it generated electric fields strong enough to accelerate electrons and protons toward each other with the centre of mass energy of 300 GeV, equivalent to the highest energy electron–proton colliding beam (HERA) built on Earth. For a strong solar flare which occurred on 14 July 2000, we computed the flux of muons which reached Earth. Our theoretical flux agrees quite nicely with the experimental data from the L3+C Collaboration at LEP through their observation of high-energy muons produced in coincidence with this huge flare [23].

In the present paper, we shall try to provide a unified picture of electroweak (EW) induced LENRs bringing out the essential physics and omitting many technical details all of which can be found in our earlier papers. Order of magnitude estimates of relevant parameters for different physical processes utilizing the electric and/or the magnetic modes will be presented to stress the feasibility of LENRs. The paper is organized as follows. In §2, general considerations revealing the basic idea behind EW induction of LENR is given. In §3, the case of metallic hydride cells where electric charge fluctuations play a major role is discussed along with estimates of the expected rates of nuclear transmutations. Estimates of the mean free path for the ultralow momentum neutrons and MeV range γ-rays are shown to be so short as to confine them to the material, i.e. they get absorbed on the surface. In §4, we consider strong electric currents flowing through thin wires to show how the collective magnetic mode energy generates a huge chemical potential in the MeV range sufficient to induce LENRs. In §5, applying the mechanism presented in §4 we show how a giant transformer – or a betatron – is generated for the solar corona leading to LENRs and, for solar flares the production of extremely high-energy particles exposing the rich structure of the Standard Model. Section 6 closes the paper with a summary of our results along with some concluding remarks and future outlook.

2. Genesis of EW-induced LENR

A free neutron can and does decay via weak interaction into a proton, an electron and a $\bar{\nu}_{\rm e}$ as

$$n \rightarrow \ p + e^- + \bar{\nu}_{\rm e}, \tag{5}$$

because the Q-value for this reaction, $Q = (M_{\rm n} - M_{\rm p} - m)c^2 \approx 0.78$ MeV is positive. On the other hand, the production of a neutron through the inverse reaction $e^- + p \rightarrow n + \nu_{\rm e}$ for electrons and protons of very low kinetic energy (generally to be found in condensed matter systems) is kinematically forbidden unless the energy of the incident (ep) system can be augmented by this Q-value. Hence, to induce LENR through a weak interaction such as

$$W + e^- + p \rightarrow n + \nu_{\rm e}, \tag{6}$$

to be kinematically allowed, we require an energy $W \geq 0.78$ MeV to be fed into the $(e^- p)$ system originally with negligible kinetic energy. Thus, our first task is to find a mechanism within a condensed matter system which can supply MeV scale energies to accelerate an electron to overcome the threshold barrier. As electrons are accelerated by an electric field through the equation $\dot{p} = (-e)\mathbf{E}$, let us assume that on a metal surface a sinusoidal electric field exists: $\mathbf{E} = \mathbf{E}_0 \cos(\Omega t)$ of frequency Ω. The average change in the momentum (Δp) is then easily obtained through $(\Delta p)^2 = e^2 \bar{E}^2/\Omega^2$. The average (squared) total energy $\bar{K}^2$ for an electron of rest mass m with an original small momentum $\mathbf{p}$ is given by

$$\bar{K}^2 = (mc^2)^2 + (c\mathbf{p})^2 + \frac{e^2 c^2 \bar{E}^2}{\Omega^2} = (c\mathbf{p})^2 + (mc^2)^2 \left[1 + \frac{\bar{E}^2}{\mathcal{E}^2}\right], \tag{7}$$

where the relevant scale of the required electric field for the neutron production is set by $\mathcal{E} = (mc/\hbar)(\hbar\Omega/e)$ [25]. For metallic hydride surfaces upon which plasma oscillations exist, typical values for the surface plasmon polariton frequencies are in the range $(\hbar\Omega/e) \approx$ (5–6) $\times 10^{-2}$ V, thereby requiring $\mathcal{E} \approx 2 \times 10^{11}$ V/m. To put it in perspective, let us recall that typical atomic electric fields are of this order of magnitude. Consider the electric field located at a Bohr radius ($a \approx 0.5$ Å) away from an isolated proton. It is given by: $\mathcal{E}(a) = (e/a^2) \approx 5 \times 10^{11}$ V/m. Coherent proton oscillations on a metallic hydride monolayer will be shown in the next section to produce electric fields and plasma frequencies of the order of magnitudes needed for neutron production. There we shall show that neutrons are born with very small momentum (ultracold) because their production is collective through a large number of protons coherently oscillating over a macroscopic region of the monolayer surface. As the nuclear absorption cross-section of ultralow momentum neutrons is extremely large, it has two desirous effects: (i) the nuclear transmutation probability is large which makes the rates substantial and (ii) the mean free path for a neutron to escape outside the metal surface is reduced to atomic distances. Hence there are no free neutrons for this process. There is, in addition, a photon shield created by mass-renormalized electrons which inhibits MeV range γ-rays to escape the surface region. We shall review it in the next section.

Let us now turn to the magnetic mode of exciting 'electron capture by proton' in a strong current-carrying wire in the steady state. For a wire of length Λ carrying a steady current I with N flowing electrons, the collective kinetic energy due to the motion of all the electrons is most simply described through the inductive energy formula

$$W = (1/2c^2)LI^2, \tag{8}$$

where $L = \eta\Lambda$, is the inductance and η (of order unity) is the inductance per unit length. If an electron is removed (by any means, such as when an electron is destroyed in a weak interaction), the change in the current is

$$\delta I = -e\left(\frac{v}{\Lambda}\right), \tag{9}$$

where v is the mean velocity of the electrons in the current. The chemical potential is then given by

$$\mu = -\frac{\partial W}{\partial N} = -\left(\frac{L}{c^2}\right) I[I(N) - I(N-1)] = \frac{e\eta I v}{c^2}. \tag{10}$$

We can write it in a more useful (system of unit independent) form using the Alfvén current $I_0 \approx 17$ kA, which was defined in eq. (4).

$$\mu = (mc^2)\eta \left(\frac{I}{I_0}\right)\left(\frac{v}{c}\right). \tag{11}$$

Thus, we see that even with a moderate $(v/c) \approx 0.1$, if currents are much larger than the Alfvén value, the chemical potential can be of the order of MeV's or higher. This is an example of how the collective magnetic kinetic energy can be distributed to accelerate a smaller number of particles with sufficient energy to produce neutrons. Further discussions about the exploding wires are given in §4.

Let us now consider the solar corona for which the magnetic flux geometry is different from that of a wire. In a wire, the magnetic field loops surround the flowing current. In the corona, there are oppositely directed currents of electrons and protons which loop around the walls of magnetic flux tubes. In a steady state magnetic flux tube which enters the solar corona out of one sun spot and returns into another sun spot without exploding, there is a substantial amount of stored magnetic energy. For a small change δI, in the current going around the vortex circumference, the change in the magnetic energy $\delta\mathcal{E}_{\rm mag}$ is given by

$$\delta\mathcal{E}_{\rm mag} = \left(\frac{1}{c}\right)\delta\Phi I. \tag{12}$$

If the length L denotes the vortex circumference, then – as described previously – for the destruction of an electron in the weak interaction, the change in the current corresponds to

$$\delta I = -e\left(\frac{v}{L}\right), \tag{13}$$

where v denotes the relative tangential velocity between the electron and the proton. Setting $\Phi = B\Delta S$ and $\delta\mathcal{E}_{\rm mag} = -W_{\rm mag}$, we obtain

$$W_{\rm mag} = (eB)\left(\frac{\Delta S}{L}\right)\left(\frac{v}{c}\right). \tag{14}$$

For a cylindrical flux tube,

$$\left(\frac{\Delta S}{L}\right) = \frac{\pi R^2}{2\pi R} = \frac{R}{2}. \tag{15}$$

For numerical estimates for the Sun, it is useful to rewrite the above as

$$W_{\rm mag} \approx (15\ {\rm GeV})\left(\frac{B}{\rm kG}\right)\left(\frac{R}{\rm km}\right)\left(\frac{v}{c}\right). \tag{16}$$

For an estimate, consider some typical values

$$R \approx 10^2 \text{ km}$$
$$B \approx 1 \text{ kG}$$
$$\frac{v}{c} \approx 10^{-2}$$
$$W_{\text{mag}} \approx 15 \text{ GeV}. \quad (17)$$

Thus, even when the flux tube does not explode, appreciable neutron production is to be expected. It should be noted that neutrons produced via weak interactions in the higher-energy regime dominated by collective magnetic effects do not necessarily have ultralow momentum.

On the other hand, for a spectacular solar flare which lasts for a time Δt, the loss of magnetic flux tube would yield a mean Faraday law acceleration voltage $\bar{V}$ around the walls given by

$$\bar{V} = \frac{\Delta\Phi}{c\Delta t}. \quad (18)$$

Inserting $\Delta\Phi = B\Delta S$ as before, where B denotes the mean magnetic field before the explosion and ΔS is the inner cross-sectional area of the flux tube, we have for the mean acceleration energy

$$e\bar{V} = (eB)\frac{\Delta S}{\Lambda}, \quad \text{where } \Lambda = c\Delta t. \quad (19)$$

For a cylindrical geometry, we can again rewrite it in a useful form

$$e\bar{V} \approx \ (30 \text{ GeV}) \left(\frac{B}{\text{kG}}\right) \left(\frac{\pi R^2}{\Lambda - \text{km}}\right). \quad (20)$$

For a coronal mass ejecting coil exploding in a time $\Delta t \approx 10^2$ s, we may estimate

$$R \approx 10^4 \text{ km}$$
$$B \approx 1 \text{ kG}$$
$$\Lambda \approx 3 \times 10^7 \text{ km}$$
$$e\bar{V} \approx 300 \text{ GeV}. \quad (21)$$

Physically, it corresponds to a colliding beam of electrons and protons with a centre of mass energy of 300 GeV. More on these matters can be seen in §5.

Having discussed the mechanisms and making ourselves familiar with the magnitudes of the parameters involved in both the collective electric and magnetic modes of the exciting neutron production, we shall devote the next three sections to a more detailed description of how it is realized in the three different physical cases: metallic hydride cells, exploding wires and the solar corona.

3. Metallic hydride cells

Even though our discussion would hold for any metallic hydride, we shall concentrate here on palladium hydrides, which are particularly suited for our purpose because on such a loaded hydride, there will be a full proton layer on the hydride surface. On this surface there will then exist coherent proton oscillations. That is, all the protons will be undergoing a collective oscillation known as the surface plasmon mode. Let us determine the size of this plasma frequency Ω and estimate the mean electric field $\bar{E} = \sqrt{\bar{\mathbf{E}}^2}$ generated on the surface.

Suppose a proton of mass $M_{\rm p}$ is embedded in a sphere with a mean electronic charge density $\rho_{\rm e} = (-e)n$. If the proton suffers a small displacement $\mathbf{u}$, then an electric field will be created

$$e\mathbf{E} = -\left(\frac{4\pi e^2 n}{3}\mathbf{u}\right) = -M_{\rm p}\Omega^2\mathbf{u}, \tag{22}$$

to satisfy Gauss' law $\operatorname{div}\mathbf{E} = 4\pi\rho_{\rm e}$. This electric field will try to push back the proton to the centre of the sphere. The equation of motion of the proton $M_{\rm p}\ddot{\mathbf{u}} = e\mathbf{E} = -\Omega^2 M_{\rm p}\mathbf{u}$ yields the oscillation. This equation also furnishes the relationship between the mean electric field and the mean proton displacement

$$e^2\bar{\mathbf{E}}^2 = \left(\frac{4\pi e^2 n}{3}\right)^2 \bar{\mathbf{u}}^2. \tag{23}$$

We can estimate the strength of the mean electric field by taking the mean electron number density at the position of the proton

$$n \approx |\psi(0)|^2 = \frac{1}{\pi a^3}. \tag{24}$$

Defining the atomic electric field at a distance a away as

$$\mathcal{E}_a = \frac{e}{a^2} \;\approx\; 5.1 \,\times\, 10^{11}\ \text{V/m}, \tag{25}$$

we obtain

$$\bar{\mathbf{E}}^2 = \mathcal{E}_a^2\left(\frac{16}{9}\right)\left(\frac{\bar{\mathbf{u}}^2}{a^2}\right). \tag{26}$$

Neutron scattering experiments on palladium hydride clearly indicate a sharply defined oscillation peak for $(\hbar\Omega/e) \approx 60$ mV, as quoted in §2. Such a collective proton motion at an infrared frequency will resonate with electronic surface plasmon oscillations leading to the local breakdown of the Born–Oppenheimer approximation [26]. They will also lead to large collective proton oscillation amplitudes. This explains the large mean proton displacement $\bar{u} \approx 2.2$ Å estimated from the neutron scattering experiments and a mean electric field estimate through eq. (26)

$$\bar{E} \approx 28.8 \times 10^{11}\ \text{V/m} \tag{27}$$

which is over 10 times larger than $\mathcal{E}$ thus proving that the plasma oscillations on metallic hydride surfaces do provide internal local electric fields more than sufficient to accelerate the surface plasmon polariton electrons in overcoming the threshold barrier.

Our description about the acceleration of an electron due to charge oscillations on the surface can be recast in a manifestly Lorentz and gauge covariant form to imply that the free electron mass m has been 'dressed up' or renormalized to a higher value $\tilde{m} = \beta m$ for the surface electrons (*vedi* eq. (7)). Once $\beta \geq \beta_0 \approx 2.53$, neutron production through weak interaction is kinematically allowed.

For these heavy surface plasmon polariton electrons $\tilde{e}^-$ with more than sufficient mass to enable the weak interaction $\tilde{e}^- + p \to n + \nu_{\rm e}$ to proceed, we turn to a rough order of magnitude estimate of this reaction rate. For this purpose, we can employ (i) the usual Fermi point-like left-handed interaction with coupling constant $G_{\rm F}$ [27], (ii) a heavy electron mass $\tilde{m} = \beta m$ and (iii) the small neutron–proton mass difference $\Delta = (M_{\rm n} - M_{\rm p}) \approx 1.3$ MeV/c^2. To make an order of magnitude estimate of the rate of this reaction, we observe that this rate which (in lowest order of perturbation theory) is proportional to $G_{\rm F}^2$ must on dimensional grounds scale with the fifth power of the electron mass. Hence, a rough dimensional analysis estimate will give

$$\Gamma(\tilde{e}^- p \to n\nu_{\rm e}) \approx \left(\frac{G_{\rm F} m^2 c}{\hbar^3}\right)^2 \left(\frac{mc^2}{\hbar}\right) \left(\frac{\tilde{m} - \Delta}{\Delta}\right)^2. \tag{28}$$

Numerically, this would imply a rate

$$\Gamma(\tilde{e}^- p \to n\nu_{\rm e}) \approx 7 \times 10^{-3} \left(\frac{\tilde{m} - \Delta}{\Delta}\right)^2 \text{ Hz}, \tag{29}$$

which in turn implies

$$\Gamma(\tilde{e}^- p \to n\nu_{\rm e}) \approx 1.2 \times 10^{-3} (\beta - \beta_0)^2 \text{ Hz}. \tag{30}$$

If we assume a surface density of 10^{16}/cm^2 (heavy electron–proton) pairs, we arrive at the following estimate for the rate of weak neutron production on a hydride surface:

$$\tilde{w}(\tilde{e}^- p \to n\nu_{\rm e}) \approx (1.2 \times 10^{13}/\text{cm}^2/\text{s})(\beta - \beta_0)^2, \tag{31}$$

which is substantial indeed.

3.1 *Theoretical computation of the rate*

Here we provide a detailed calculation of the neutron production cross-section from a proton and an electron of mass ($\tilde{m}$) through the Fermi interaction. The Feynman amplitude for the reaction

$$\tilde{e}(p; s_{\rm p}) + p(P_{\rm p}; S_{\rm p}) \to n(P_n; S_n) + \nu(k; \lambda) \tag{32}$$

in the Born approximation reads (in natural units $\hbar = c = 1$)

$$\mathcal{M} = \frac{G_F}{\sqrt{2}} H^\mu L_\mu, \tag{33}$$

where

$$H^\mu = \bar{U}(P_n; S_n)\gamma^\mu (C_v - C_A\gamma_5) U(P_p; S_p) \tag{34}$$

and

$$L_\mu = \bar{u}(k; \lambda)\gamma_\mu (1 - \gamma_5) u(p; s_p). \tag{35}$$

In the following, we shall for simplicity assume $C_v = C_A = 1$. Then, the spin-averaged cross-section reads

$$\sigma(\tilde{e} + p \to\ n + \nu) = \left[\frac{2G_F^2}{\pi}\right] \left[\frac{p_f}{p_i}\right] \left[\frac{(P_n \cdot k)(P_p \cdot p)}{s}\right], \tag{36}$$

where $s = (P_p + p)^2 = (P_n + k)^2$ and $p_{i,f}$ denote the magnitudes of the initial and final CM 3-momenta. This reaction is physically realizable once $\tilde{m} > \Delta$ and then it is exothermic. Hence, the cross-section in eq. (36) diverges as $1/p_i$ as $p_i \to 0$ but the rate of the reaction which is proportional to $v_i\sigma$ is finite. Using eq. (36), we can compute $v_i\sigma$ in the limit of vanishing relative velocity v_i as

$$v_i\sigma(\tilde{e} + p \to n + \nu) \approx \left[\frac{2G_F^2}{\pi}\right] (\tilde{m} - \Delta)^2. \tag{37}$$

This explicit calculation justifies the order of magnitude estimates made in eq. (28) *et seq* and in particular the claim that the neutron production rate is indeed proportional to $(\tilde{m} - \Delta)^2$.

3.2 *Nuclear transmutations through the produced neutrons*

The neutrons produced above will be of ultralow momentum because their production is through the collectively oscillating protons acting in tandem from a patch on the surface. The wavelength of the neutrons can be estimated to be about $\lambda \approx 10^{-3}$ cm. Such long wavelength neutrons will get absorbed with an extremely high probability by the nuclei on the surface because the neutron absorption cross-section will be very large. This can be seen by computing the total neutron cross-section through the optical theorem which relates it to the forward elastic (n-Nucleus) amplitude

$$\sigma_T(n + \text{Nucleus} \to \text{anything}) = \left(\frac{4\pi}{k}\right) \Im m f(k, \mathbf{0}). \tag{38}$$

Let $f(k, \mathbf{0}) = a + ib$, with $b \approx 1$ Fermi. Then, we obtain

$$\sigma_T(n + \text{Nucleus} \to \text{anything}) = 2\lambda b \approx 2 \times 10^8 \text{ b}, \tag{39}$$

a very large value. This not only shows a hefty rate for the production of nuclear transmutations through a rapid absorption of neutrons but it also shows that the mean free path of a neutron Λ is of the order of a few atomic distances. In fact, given the density of neutron absorbers $n_{\rm abs} \approx 10^{22}/{\rm cm}^3$, we may estimate

$$\Lambda^{-1} = n_{\rm abs}\sigma_{\rm T}; \ \Lambda \approx 50 \ \text{Å}. \tag{40}$$

Hence, practically all produced neutrons will get absorbed with essentially a zero probability of finding a free neutron.

The observed electromagnetic radiation from the surface heavy electrons will be confined essentially to low-energy photons reaching up to soft X-rays with practically no MeV range photon's being radiated because the mean free path of the produced γ-rays in the few MeV range would be very short, about a few Angstroms.

To recapitulate: The surface charge oscillation plasmons provide enough collective energy for the production of heavy mass electrons which in turn lead to the production of low-momentum neutrons. Such neutrons get readily absorbed and their production dynamics produces their own neutrons and built-in γ-ray 'shields'. The observable end products would just be the nuclear transmutations triggered by the absorbed neutrons. A plethora of nuclear reactions are thereby possible. One such complete nuclear chain cycle with a high Q-value is as follows. Let us assume that the surface is coated with lithium. Successive absorption of neutrons by lithium will produce ${}^4_2$He:

$$\begin{aligned}
&{}^6_3\mathrm{Li} + n \to {}^7_3\mathrm{Li} \\
&{}^7_3\mathrm{Li} + n \to {}^8_3\mathrm{Li} \\
&{}^8_3\mathrm{Li} \to {}^8_4\mathrm{Be} + e^- + \bar{\nu}_{\rm e} \\
&{}^8_4\mathrm{Be} \to {}^4_2\mathrm{He} + {}^4_2\mathrm{He}.
\end{aligned} \tag{41}$$

The heat produced by the above reaction is quite high: $Q[{}^6_3\mathrm{Li} + 2n \to 2\,{}^4_2\mathrm{He} + e^- + \bar{\nu}_{\rm e}] \approx 26.9$ MeV.

On the other hand, ${}^4_2$He can successively absorb neutrons and, through the formation of intermediate halo nuclei, reproduce lithium

$$\begin{aligned}
&{}^4_2\mathrm{He} + n \to {}^5_2\mathrm{He} \\
&{}^5_2\mathrm{He} + n \to {}^6_2\mathrm{He} \\
&{}^6_2\mathrm{He} \to {}^6_3\mathrm{Li} + e^- + \bar{\nu}_{\rm e}.
\end{aligned} \tag{42}$$

The heat from the reaction in eq. (42) is $Q[{}^4_2\mathrm{He} + 2n \to {}^6_3\mathrm{Li} + e^- + \bar{\nu}_{\rm e}] \approx 2.95$ MeV. The complete nuclear cycle as described in eqs (41) and (42) taken together releases a substantial total heat through nuclear transmutations. Other lithium-initiated processes produce both ^{4}He and ^{3}He.

3.3 *Isospin precession*

Consider the nuclear mass formula, i.e., tha mass $M(Z, A)$ of a nucleus consisting of Z protons and $N = (A - Z)$ neutrons where A is the total number of nucleons in the nucleus. It reads

$$M(Z,A)c^2 = AM_{\rm n}c^2 + Z(M_{\rm p} - M_{\rm n})c^2 - B(Z,A), \tag{43}$$

where the binding energy is given by

$$-B = -\epsilon_1 A + \epsilon_2 A^{2/3} + \epsilon_3 \left(\frac{Z^2}{A^{1/3}}\right) + \epsilon_4 \frac{(A-2Z)^2}{A} + \epsilon_5 \frac{\lambda}{A^{3/4}}, \tag{44}$$

where $\epsilon_1 = 15.75$ MeV, $\epsilon_2 = 17.8$ MeV, $\epsilon_3 = 0.71$ MeV, $\epsilon_4 = 23.7$ MeV and $\epsilon_5 = 34$ MeV. The parameter λ assumes the value of (i) +1 for odd–odd nuclei; (ii) 0 for even–odd and (iii) −1 for even–even nuclei.

One can therefore, define a nuclear voltage $\Phi(Z,A)$ as

$$e\Phi(Z,A) = c^2 \frac{\partial M(Z,A)}{\partial Z}. \tag{45}$$

Stability of a nucleus will hence be achieved for $\Phi(Z^*,A) = 0$, with Z^* given by

$$Z^* = \frac{A}{2 + (\epsilon_3/2\epsilon_4)A^{2/3}} \approx \frac{A}{2 + 0.015A^{2/3}}. \tag{46}$$

Thus, stable nuclei lie on a plot of $N^* = (A - Z^*)$ vs. Z^* rather than on the symmetry axis $N = Z$.

For $Z < 50$, stable nuclei would arise through the weak β^- decay from the neutron-rich unstable nuclei and arise from the weak β^+ decay from the proton-rich unstable nuclei. Hence, low (negative) voltage Φ^- yields β^- decay and high (positive) voltage Φ^+ yields β^+ decay.

Now, just as the Earth's magnetic moment exhibits a precession or as the electron magnetic moments exhibit a precession, we expect a nuclear isospin rotation to exhibit a weak precession due to the nuclear voltage Φ.

Consider the nuclear isospin vector $\mathbf{T} = (T_1, T_2, T_3)$ satisfying the algebra

$$[T_i, T_j] = i\epsilon_{ijk}T_k, \tag{47}$$

with

$$\mathbf{T}\cdot\mathbf{T} = T(T+1), \quad T_3 = \frac{(Z-N)}{2} = Z - \frac{1}{2}A. \tag{48}$$

Hence, with our phenomenological 'Hamiltonian' $M(Z,A)c^2$, we expect a precession of the isospin vector

$$\frac{\partial \mathbf{T}}{\partial \tau} = \frac{i}{\hbar}[M(Z,A)c^2, \mathbf{T}] = \boldsymbol{\Omega} \times \mathbf{T}, \tag{49}$$

where

$$\hbar\boldsymbol{\Omega} = (0, 0, e\Phi), \tag{50}$$

with the frequency of precession given by $\Omega = (e\Phi/\hbar)$.

Defining as usual $T_\pm = T_1 \pm iT_2$, the above implies

$$T_{\pm}(\tau) = T_{\pm}(0)\mathrm{e}^{\pm i\Omega\tau}. \tag{51}$$

This isospin rotation describes a 'virtual' W decay into electron and neutrino excitations with $T_{\pm}$ yielding $W^{\mp} \to \beta^{\mp} + \nu^{\pm}m$, where $\nu^{+} = \bar{\nu}$ and $\nu^{-} = \nu$.

Such isospin wave excitations can be induced by long-ranged electrodynamics. The explicit coupling between isospin wave charges are given by

$$\bar{\Phi}_a = \bar{\Phi}_a^{\mathrm{C}} + \bar{\Phi}_a^{\mathrm{D}}, \tag{52}$$

with the Coulomb term given by

$$\bar{\Phi}_a^{\mathrm{C}} = e\left\langle \sum_{b\neq a} \frac{Z_b}{|\mathbf{R_b} - \mathbf{R_a}|} - \sum_j \frac{1}{|\mathbf{R_a} - \mathbf{r_j}|} \right\rangle, \tag{53}$$

and the velocity-dependent Darwin term (discussed in detail in the next section) reads

$$\bar{\Phi}_a^{\mathrm{D}} = -e\left\langle \sum_{b\neq a} \frac{Z_b(\mathbf{v}_a \cdot \mathbf{v}_b + \mathbf{v}_a \cdot n_{ab}\mathbf{v}_b \cdot n_{ab})}{c^2|\mathbf{R_b} - \mathbf{R_a}|} - \sum_j \frac{(\mathbf{v}_a \cdot \mathbf{v}_j + \mathbf{v}_a \cdot n_{aj}\mathbf{v}_j \cdot n_{aj}}{c^2|\mathbf{R_a} - \mathbf{r_j}|} \right\rangle. \tag{54}$$

The detailed dynamics and the technical challenge of such exciting long-range isospin waves will be discussed elsewhere.

4. Exploding wires

In §2, we have outlined our explanation of nuclear transmutations and fast neutrons which have been observed to emerge when large electrical current pulses passing through the wire filaments are induced to explode. If a strong current pulse, large on the scale of I_0, defined in eq. (4), passes through a thin wire filament, then the magnetic field which loops around the current, exerts a very large Maxwell pressure on surface area elements, compressing, twisting and pushing into the wire. When the magnetic Maxwell pressure grows beyond the tensile strength of the wire material at the hot filament temperature, the wire first expands, then begins to melt and finally disintegrates. There now exist slow-motion pictures which verify that indeed the wire expands, melts and disintegrates. All of this is readily understood. If the heating rate is sufficiently fast, then the hot wire may emit thermal radiation at a very high noise temperature. The thermal radiation from the exploding tungsten filaments exhibits X-ray frequencies indicating very high electron kinetic energies within the filament. Due to the electron kinetic pressure, the wire diameter starts to increase yielding a filament-dense gas phase but still with some liquid droplets. The final explosive product consists of a hot plasma colloid containing some small dust particles of the original wire material. These products cool off into a gas and some smoke as is usual for explosions.

As discussed in §2, we want to understand how LENRs can be initiated in an exploding wire current pulse with a strong current (with its peak value substantially higher than I_0) produced by a capacitor discharge with an initial voltage of only 30 keV [28,29]. We also want to understand, by contrast, why when Rutherford had fired a much higher energy 100 keV but dilute beam of electrons into a tungsten target he did not observe any nuclear reactions [30].

A typical electron in the current with a mean kinetic energy of 15 keV would have an average speed $(v/c) \approx 0.25$. On the other hand, even for such low mean speed, the chemical potential given in eq. (11), for $(I/I_0) \approx 200$ becomes large

$$\mu \approx (mc^2)(200)(0.25) = 25 \text{ MeV}, \tag{55}$$

comfortably sufficient for an electron to induce a weak interaction LENR. Overall energy conservation will of course require that only a certain fraction of about (15 keV/25 MeV) $= 6 \times 10^{-4}$ of the total number of electrons in the current will be kinematically allowed to undergo weak interactions.

Let us now briefly discuss why Rutherford with his much higher energy – but dilute – beam of electrons did not observe any nuclear reactions. The reason is rather simple. In the vacuum, there is a mutual Coulomb repulsion between the electrons in the beam which compensates the mutual Amperian current attraction. In the exploding wire filament, on the other hand, the repulsive Coulomb part is screened by the background positive charge but leaves intact the Amperian current attraction thereby allowing the possibility of nuclear reactions.

In the following subsection, we give a theoretical description based on the Darwin Lagrangian to justify the general ideas discussed above.

4.1 *Darwin electrodynamics*

Let us go back to Darwin. In 1920, Darwin [31] constructed an effective action for electrodynamics – valid to order c^{-2} – in which the near field vector potentials were eliminated in favour of instantaneous but velocity-dependent forces. To a sufficient degree of accuracy, for our purposes, one may write a Lagrangian for a system of point particles with positions $\mathsf{R} = (\mathbf{r}_1, \ldots, \mathbf{r}_N)$, velocities $\mathsf{V} = (\mathbf{v}_1, \ldots, \mathbf{v}_N)$, masses $(m_1, \ldots, m_N)$ and charges $(e_1, \ldots, e_N)$. It is

$$\begin{aligned}
L(\mathsf{V},\mathsf{R}) &= K(\mathsf{V},\mathsf{R}) - U(\mathsf{R}),\\
K(\mathsf{V},\mathsf{R}) &= \frac{1}{2}\sum_{a=1}^{N} m_a|\mathbf{v}_a|^2 + \tilde{K}(\mathsf{V},\mathsf{R}),\\
\tilde{K}(\mathsf{V},\mathsf{R}) &= \sum_{a<b}^{N} \frac{e_a e_b[\mathbf{v}_a\cdot\mathbf{v}_b + (\mathbf{v}_a\cdot\mathbf{n}_{ab})(\mathbf{v}_b\cdot\mathbf{n}_{ab})]}{2c^2 r_{ab}},\\
U(\mathsf{R}) &= \sum_{a<b}^{N} \frac{e_a e_b}{r_{ab}},
\end{aligned} \tag{56}$$

where $\mathbf{r}_{ab} = \mathbf{r}_a - \mathbf{r}_b$ and $\mathbf{n}_{ab} = \mathbf{r}_{ab}/r_{ab}$. In eqs (56), $\tilde{K}(\mathsf{V},\mathsf{R})$ represents the magnetic near-field interaction between slowly moving charges when viewed as an effective contribution to the kinetic energy and $U(\mathbf{R})$ represents the Coulomb potential energy. The total energy $\mathcal{E}$ obtained from the Lagrangian eqs (56) is by the usual rule the sum of the generalized kinetic and potential energy

$$\mathcal{E} = \left(\sum_{a=1}^{N} \mathbf{v}_a \cdot \frac{\partial L}{\partial \mathbf{v}_a} \right) - L,$$
$$\mathcal{E}(\mathsf{V}, \mathsf{R}) = K(\mathsf{V}, \mathsf{R}) + U(\mathsf{R}). \tag{57}$$

In deriving eqs (57), we have used the fact that the kinetic energy is homogeneous of degree two, i.e. $K(\lambda \mathsf{V},\ \mathsf{R}) = \lambda^2 K(\mathsf{V}, \mathsf{R})$ and so by virtue of Euler's theorem $\sum_a \mathbf{v}_a \cdot (\partial K / \partial \mathbf{v}_a) = 2K$. Note that the generalized kinetic energy K is no longer a sum of the single particle kinetic energies. The magnetic interaction describes a two-body collective contribution to kinetic energy K thereby the conserved total energy $\mathcal{E}$.

Let us consider the motion of charged particles flowing through a straight thin wire of self-inductance L and length Λ directed along a unit vector $\mathbf{n}$. For the purpose at hand, it is useful to concentrate on the Darwin kinetic energy adapted to the inductive circuit element form [32] written in terms of current I and inductance L:

$$K_{\text{Darwin}} = \sum_a \frac{1}{2} m_a v_a^2 + \frac{1}{2c^2} L I^2, \tag{58}$$

where the current is given by

$$I = \left(\frac{1}{\Lambda} \right) \sum_a e_a (\mathbf{n} \cdot \mathbf{v}_a), \tag{59}$$

so that

$$K_{\text{Darwin}} = \sum_a \frac{1}{2} m_a v_a^2 + \left(\frac{\eta}{2c^2 \Lambda} \right) \sum_{b \neq a} e_a e_b (\mathbf{n} \cdot \mathbf{v}_a)(\mathbf{n} \cdot \mathbf{v}_b). \tag{60}$$

Here, $\eta = L/\Lambda$ is the (dimensionless) inductance per unit length of the wire.

As in the earlier case, eq. (60) shows explicitly that the kinetic energy of a given charged particle depends on the velocities of the other charged particles in the current flow. We may compute the change in the momentum $\mathbf{p}$ and the energy W of a given electron through the collective effect of the entire current in the following way. If the current changes rapidly, it will induce a Faraday law voltage across an inductive current element. This voltage (per unit length) is the induced electric field $\mathbf{E}$, whose magnitude is given by

$$E = \left(\frac{\eta}{c^2} \right) \left(\frac{\mathrm{d}I}{\mathrm{d}t} \right). \tag{61}$$

The evident equation of motion of an electron (under this electric field) is $\dot{\mathbf{p}} = e\mathbf{E}$ and the power delivered to a moving electron by a changing current is given by

$$\frac{\mathrm{d}W}{\mathrm{d}t} = evE = -\eta(mc^2)\left(\frac{1}{I_0}\frac{\mathrm{d}I}{\mathrm{d}t}\right)\frac{v}{c}, \tag{62}$$

where we have introduced the Alfvén current through $I_0/c = R_{\mathrm{vac}}I_0/4\pi = (mc^2)/e$, so that $I_0 \approx 1.704509 \times 10^4$ A. I_0 is the current corresponding to the electron rest mass and thus currents of this order or larger are expected to lead to interesting physics at the MeV (or nuclear physics) scale. This provides a sound theoretical basis for the applications described above and is also central to the solar corona and solar flares described below.

5. Solar corona and flares

As stated in §2, oppositely directed Amperian currents of electrons and protons loop around the walls of a magnetic flux tube which exits out of one sun spot into the solar corona to enter back into another sun spot. The magnetic flux tube is held up by magnetic buoyancy. We consider here the dynamics of how very energetic particles are produced in the solar corona and how they induce nuclear reactions well beyond the solar photosphere. Our explanation, centred around Faraday's law, produces the notion of a solar accelerator very similar to a betatron [33,34]. A betatron is a step-up transformer whose secondary coil is a toroidal ring of particles circulating around a time-varying Faraday flux tube.

We can view the solar flux tube to act as a step-up transformer which passes some circulating particle kinetic energy from the solar photosphere outward to other circulating particles in the solar corona. The circulating currents within the photosphere are to be considered as a net current $I_{\mathcal{P}}$ around a primary coil and the circulating currents high up in the corona as a net current $I_{\mathcal{S}}$. If $K_{\mathcal{P}}$ and $K_{\mathcal{S}}$ represent the kinetic energies, respectively, in the primary and the secondary coils, the step-up transformer power equation reads

$$|\dot{K}_{\mathcal{P}}| = |V_{\mathcal{P}}I_{\mathcal{P}}| = |V_{\mathcal{S}}I_{\mathcal{S}}| = |\dot{K}_{\mathcal{S}}|, \tag{63}$$

where $V_{\mathcal{P}}$ and $V_{\mathcal{S}}$ represent the voltages across the primary and the secondary coils, respectively. The total kinetic energy transfer reads

$$\Delta K_{\mathcal{P}} = \int(\mathrm{d}t)|V_{\mathcal{P}}I_{\mathcal{P}}| = \int(\mathrm{d}t)|V_{\mathcal{S}}I_{\mathcal{S}}| = \Delta K_{\mathcal{S}}. \tag{64}$$

In essence, what the step-up transformer mechanism does is to transfer the kinetic energy distributed amongst a very large number of charged particles in the photosphere – via the magnetic flux tube – into a distant much smaller number of charged particles located in the solar corona, so that a small accelerating voltage in the primary coil produces a large accelerating voltage in the secondary coil. The transfer of kinetic energy is collective from a larger group of particles into a smaller group of particles resulting in the kinetic energy per charged particle of the dilute gas in the corona becoming higher than the kinetic energy per particle of the more dense fluid in the photosphere.

We can convert the above into a temperature–kinetic energy relationship by saying that the temperature of the dilute corona will be much higher than the temperature of the more dense fluid photosphere. If and when the kinetic energy of the circulating currents in a part of the floating flux tube becomes sufficiently high, the flux tube would become unstable and explode into a solar flare which may be accompanied by a coronal mass ejection. There is a rapid conversion of the magnetic energy into charged particle kinetic energy. These high-energy products from the explosion initiate nuclear as well as elementary particle interactions, some of which have been detected in laboratories.

Recent NASA and ESA pictures show that the surface of the Sun is covered by a carpet-like interwoven mesh of magnetic flux tubes of smaller dimensions. Some of these smaller structures possess enough magnetic energy to lead to LENRs through a continual conversion of their energy into particle kinetic energy. Occurrence of such nuclear processes in a roughly steady state would account for the solar corona remaining much hotter than the photosphere. Needless to say that our picture belies the notion that all nuclear reactions are contained within the core of the Sun. On the contrary, it provides strong theoretical support for experimental anomalies [35,36] such as short-lived isotopes [37–39] that have been observed in the spectra of stars having unusually high average magnetic fields.

For the transformer mechanism to be fully operational in the corona, the coronal electrical conductivity must not be too large. Useful experimental bounds on an upper limit to this conductivity can be obtained through its effect on measurements of gravitational bending of light near the Sun as it traverses the solar corona. Successful measurements of the gravitational bending of electromagnetic waves with frequencies in the visible and all the way down to the high end of the radiospectrum are legion. These experiments provide a direct proof that any coronal conductivity disturbance on the expected gravitational bending of electromagnetic waves for frequencies down to 12.5 GHz must be negligible. Error estimates from even lower frequency radiowave probes used for gravitational bending [40,41] put the coronal conductivity in the megahertz range. For comparison, we note that the typical conductivity of a good metal would be more than ten orders of magnitude higher [45]. The corona is close to being an insulator and eons away from being a metal and their is no impediment toward sustaining electrical fields within it [42–44]. Thus, our proposed transformer mechanism and its subsequent predictions for the corona remain intact.

The spectacular solar flare, which occurred on 14 July 2000 and the measurement of the excess muon flux associated with this flare by the CERN L3+C group [23] offered a unique opportunity to infer that protons of energies greater than 40 GeV were produced in the solar corona. Likewise, the BAKSAN underground muon measurements [47] provided evidence for protons of energies greater than 500 GeV in the solar flare of 29 September 1989. The very existence of primary protons in this high-energy range provides strong evidence for the numbers provided in eq. (21). Hence, for large solar flares in the corona, electrons and protons must have been accelerated well beyond anything contemplated by the standard solar model. This in turn provides the most compelling evidence for the presence of large-scale electric fields and the transformer or betatron mechanism because we do not know of any other process that could accelerate charged particles to beyond even a few GeV, let alone hundreds of GeVs.

5.1 *Total rate of positron production in a flare*

Here we estimate the total rate of positrons produced in a solar flare through the reaction $e^- + p \rightarrow e^+e^- + X$. The rate of production of e^+e^- pairs is equal to the rate of production of $\mu^+\mu^-$ pairs. After a while, however, all the muons will decay and from each muon (outside the corona) we shall get one electron (or one positron). The cross-section for $\mu^+\mu^-$ pair production at HERA is 6.5 pb. We must also include the $\tau^+\tau^-$ channel. HERA data for tau-pair production lead to a cross-section of $(13.6 \pm 4.4 \pm 3.7)$ pb [46].

As a μ^+ only decays leptonically, each μ^+ will decay into an e^+. On the other hand, the branching ratio of a τ^+ decay into (i) a μ^+ is 17.36% and into (ii) an e^+ is 19.59%. Because the μ^+ would decay into an e^+, the combined branching ratio from purely leptonic channels

$$\mathrm{BR}(\tau^+_{\mathrm{leptonic}} \rightarrow\ e^+ + X) = 35.2\%. \tag{65}$$

The rest of the τ^+ decay channels contain a $\bar{\nu}_\tau$ plus hadrons with a total unit positive charge, which means either a π^+ or a K^+ along with neutrals or a couple of $\pi^+\pi^-$. Eventually, every hadronic decay channel contains at least one e^+ (some channels even have two e^+'s). Conservatively, we may take that each τ would produce one positron. Taking the central value, the cross-section for e^+ production through the tau's would be

$$\sigma(e^-p \rightarrow\ \tau^+X \rightarrow\ e^+X') = 13.6\ \mathrm{pb}. \tag{66}$$

Hence, the total e^+ production cross-section (due to all the three leptons) becomes

$$\sigma(e^-p \rightarrow\ e^+X) = (6.5 + 6.5 + 13.6)\ \mathrm{pb} = 26.6\ \mathrm{pb}. \tag{67}$$

Of course, some of these positrons might get annihilated before getting out from the corona. Let us estimate the annihilation rate of a positron within the corona (or, the lifetime of a positron in the corona). At high energies, the e^+e^- annihilation cross-section is given by

$$\sigma(e^+e^- \rightarrow\ \gamma\gamma; E_+) = \left(\frac{\pi\alpha^2}{m_{\mathrm{e}}E_+}\right)[\ln(2E_+/m_{\mathrm{e}}) - 1], \tag{68}$$

where E_+ is the 'laboratory' energy of the positron. For illustrative purposes, let us take a positron of energy $E_+\ =\ 10$ GeV. Then

$$\sigma(e^+e^- \rightarrow\ \gamma\gamma; E_+ = 10\ \mathrm{GeV}) \approx\ 0.115\ \mathrm{mb}. \tag{69}$$

Assuming the density of electrons (a few solar radii away) to be $\rho_{\mathrm{e}} = 2.78 \times 10^7/\mathrm{cm}^3$ and we know that it drops considerably as one gets further away, but for now let us assume it to be a constant, we can estimate the rate of a positron annihilating before exiting the corona to be

$$\mathrm{Rate}(e^+\mathrm{Annihil}) = \rho_{\mathrm{e}} v \sigma_{\mathrm{ann}} \approx\ 9.4 \times 10^{-11}\ \mathrm{Hz}. \tag{70}$$

Hence, the probability that a positron is annihilated before reaching the Earth's atmosphere in about 8 min, would be about 4.5×10^{-8} and thus negligible.

The rate of positron production $\dot{N}(E_{\rm CM})$ can be written as

$$\dot{N}(E_{\rm CM}) = V n_{\rm e} n_{\rm p} v_{\rm rel} \sigma(e^- \ + \ p \ \to \ e^+ + \ X)(E_{\rm CM}), \tag{71}$$

where V denotes the fiducial interaction volume, $n_{\rm e}$, $n_{\rm p}$ denote the electron and proton densities respectively, $v_{\rm rel} = |\mathbf{v}_{\rm e} - \mathbf{v}_{\rm p}|$ denotes their relative velocities, and $\sigma(e^- + p \to e^+ + X)(E_{\rm CM})$ is the inclusive positron production cross-section at the CM energy $E_{\rm CM}$. For $E_{\rm CM} = 300$ GeV, the value of this cross-section has been deduced in eq. (67) to be 26.6 pb. Also, for the solar flare considered, $V = 9.43 \times 10^{30}$ cm^3. As for the (common electron and proton) density, we shall take a conservatively small value $n_{\rm e} = n_{\rm p} = 2.78 \times 10^7/{\rm cm}^3$ [41]. Inserting these values in eq. (71), we obtain

$$\dot{N}(300 \text{ GeV}) \approx \ 11.2 \times 10^{21}/\text{s}. \tag{72}$$

Under the simplifying assumption that the positron production is isotropic, the differential positron flux before reaching the Earth's atmosphere is given by

$$F(e^+) = \frac{\dot{N}(300 \text{ GeV})}{4\pi (1.5 \times 10^{11} \text{ m})^2} \approx \ 0.04/\text{m}^2\text{-s-sr}. \tag{73}$$

This should be compared with the overall positron flux estimate for all cosmic rays (integrated over positron energies >8.5 GeV) which is about $0.12/\text{m}^2$-s-sr. Thus, our acceleration mechanism is not only capable of accelerating electrons and protons in a solar flare to hundreds of GeV but it also yields a high-energy positron flux which is a substantial fraction of the overall cosmic ray positron flux. We are unaware of any similar theoretical estimate in the literature.

The estimate of positron flux given in eq. (73) is the one directly generated through lepton pairs. There will be additional positrons (on an average with smaller energies) from the rest of the process $\gamma^*(Q) + p \to X$, where $\gamma^*(Q)$ stands for a virtual photon of mass Q. The collection X of hadrons contains a proton or a neutron along with charged and neutral pions and kaons. We can estimate this 'extra' flux of positrons by relating it to that of the protons in the following way. For each proton, we would have a certain number of charged ($\pi^+\pi^-$; K^+K^-) and neutral ($\pi^0\pi^0$; $K^0\bar{K}^0$) pairs. As we expect an equal number of charged and neutral pairs of mesons and each positively charged meson will eventually produce a positron, the number of produced positrons associated with each proton will be approximately one half of the charged particle multiplicity $\langle n \rangle_{\rm charged}$. Also, the average energy of a produced positron will be $(E/4\langle n \rangle_{\rm charged})$ produced in association with a proton of energy E.

5.2 *Total proton flux estimate for the 14 July 2000 solar flare*

As mentioned earlier, the L3 + C Collaboration measured the muon flux from 14 July 2000 solar flare arrived at their detector. Through this measurement, they were able to estimate the primary proton flux for protons with energies greater than 40 GeV. In this section we compare their value with an estimate of the overall cosmic ray flux of protons with energies greater than 40 GeV.

Let us estimate the integrated cosmic flux of primary protons (before reaching the atmosphere). From cosmic rays section of PDG, we find (after performing an integration with a power-law exponent $\alpha = 3$)

$$F_{\text{cosmic protons}}(E > 40\ \text{GeV}) \approx \frac{6 \times 10^{-3}}{\text{cm}^2\text{-s-sr}}, \tag{74}$$

to be compared with the L3 Collaboration estimate of the primary proton flux from the giant solar flare of 14 July 2000

$$F_{L3\ \text{flare protons}}(E > 40\ \text{GeV}) \approx \frac{2.6 \times 10^{-3}}{\text{cm}^2\text{-s-sr}}, \tag{75}$$

which is a significant fraction of the total cosmic ray proton flux. It is in reasonable agreement with the neutron monitors which report a fraction ranging between 0.2 and 0.6 as the increase in the number of observed particles for the same flare as compared to the background cosmic ray particle yields.

The above result is quite significant in that our proposed mechanism of acceleration is unique in predicting primary protons from a solar flare in this very high-energy range.

6. Summary and concluding remarks

We can summarize by saying that three seemingly diverse physical phenomena, viz., metallic hydride cells, exploding wires and the solar corona, do have a unifying theme. Under appropriate conditions which we have now well delineated, in all these processes electromagnetic energy gets collectively harnessed to provide enough kinetic energy to a certain fraction of the electrons to combine with protons (or any other ions present) and produce neutrons through weak interactions. The produced neutrons then combine with other nuclei to induce low-energy nuclear reactions and transmutations. Lest it escape notice let us remind the reader that all three interactions of the Standard Model (electromagnetic, weak and nuclear) are essential for an understanding of these phenomena. Collective effects, but no new physics for the acceleration of electrons beyond the Standard Model needs to be invoked. We have seen, however, that certain paradigm shifts are necessary. On the surface of a metallic hydride cell with surface plasmon polariton modes, protons collectively oscillate along with the electrons. Hence, the Born–Oppenheimer approximation (which assumes that the proton is rigidly fixed) breaks down and should not be employed. Similarly, in the solar corona, the electronic density and the electrical conductivity are sufficiently low. Hence there is not much charge screening of the electric fields involved. Strong electric fields generated by time-dependent magnetic fields through Farday's laws are sustained in the corona and the betatron (or transformer) mechanism remains functional. Were it not so, electrons and protons could not have been accelerated to hundreds of GeV's and there would have been no production of high-energy muons, certainly not copious enough to have reached Earth in sufficient numbers to have been observed by the L3 + C collaboration at LEP [23] or by the BAKSAN underground laboratory [47]. We are unaware of any other alternative scheme for obtaining this result. The betatron mechanism also naturally explains a variety of observed experimental results such as unexpected nuclear transmutations and high-energy cosmic rays from the exterior of the Sun or any other astronomical object endowed with strong enough magnetic activity such as active galactic nuclei. Also, our estimate of the muons detected at CERN are consistent with the CERN data on the Solar flare of 14 July 2000.

The analysis presented in this paper leads us to conclude that realistic possibilities exist for designing LENR devices capable of producing 'green energy', that is, production of excess heat at low cost without lethal nuclear waste, dangerous γ-rays or unwanted neutrons. The necessary tools and the essential theoretical know-how to manufacture such devices appear to be well within the reach of the technology available now. Vigorous efforts must now be made to develop such devices whose functionality requires all three interactions of the Standard Model acting in concert.

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