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Leptonic Electroweak Model with No Ad Hoc Parity-Breaking Prescriptions and No Higgs Particles

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The leptonic sector of the standard model is revisited in the light of a recent strictly covariant, generalized quantum field formalism that naturally accounts for the "maximal *P*-violation" effect without needing to postulate any true failure of *P* mirror symmetry. In full accordance with neutrino–antineutrino phenomenology, such an approach strictly predicts the nonexistence of right-handed (massless) neutrinos and left-handed (massless) antineutrinos: it spontaneously leads to a universal theory of a merely left-handed massless fermion and a merely right-handed corresponding antifermion, where such particles may now be *P*-symmetrically interpreted as just two pseudoscalar-charge conjugated objects being the ordinary mirror image of each other. On these grounds, a weak isospin of pseudoscalar components is introduced, which can alternately yield a double—either $SU(2)_L$ (leptonic) or $SU(2)_R$ (antileptonic)—variety of SU(2) representations, each with its own weak-hypercharge variety. Hence two equally allowable $SU(2) \otimes U(1)$ massless gauge schemes are covariantly built, the former including (right-handed) antileptons as $SU(2)_L$-singlets and the latter including (left-handed) leptons as $SU(2)_R$-singlets. In either alternative scheme, the "charged" lepton and antilepton can but bear as yet an electric charge quite indefinite in sign: they can really be made two electric charge conjugated eigenstates (with opposite eigenvalues) only by acquiring mass and turning into Dirac particles. Herein the charged-lepton mass appearance no longer looks like just an external "accident" connected with a hypothetical coupling to a Higgs boson; it rather becomes an essential internal requirement for the generation of the leptonic electromagnetic gauge coupling. A suitable effective mechanism of mass production is shown to be provided by simply introducing a net "vacuum" Higgs isodoublet field merely including the three Goldstone freedom degrees and not including the Higgs-boson excitation: such a field still has a nonzero vacuum expectation value and still generates SSB and Weinberg mixing, though with the only further (internal) effect of giving mass to both the intermediate vector bosons and the charged lepton (antilepton). Two equivalent covariant electroweak final schemes for either leptons or antileptons are thus obtained which contain no Higgs-boson couplings and can all the same reproduce the

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well-established standard results. Either scheme may be said not to be *P*-symmetric but for the presence of the neutral weak current coupling, as a mere outcome of the mixed (and not purely pseudoscalar) nature of the "electroweak" effective charge therein involved.

1. INTRODUCTION: A GENERALIZED, STRICTLY COVARIANT, QUANTUM FIELD FORMALISM FOR FERMIONS AND ANTIFERMIONS

Recently I have worked out a new approach to the fermionic sector of QFT[1] which can naturally account for the maximal *P*-violation effect[2,3] without appealing to ad hoc external prescriptions[4] such as the *V–A* scheme[5–7] and the neutrino two-component scheme.[8–11] Here I sum up the distinctive lines of that approach, so as to make the present paper as self-contained as possible.

The customary fermion field formalism, based on the Dirac equation and the "hole" interpretation for negative frequencies, is obviously unable to provide a one-particle relativistic description: it deals with a Fock space being the sum of two pure positive-energy Fock spaces—pertaining to "particles" and "holes", respectively—which are mapped onto each other by charge conjugation, or "particle" $\rightleftarrows$ "hole" conjugation. A strictly covariant, one-particle description is made viable once the Stüeckelberg–Feynman improved approach to the negative-energy problem[12] is used: the motion of a "hole" can then be reinterpreted as a motion, *backward* in time, of a negative-energy "particle," and the composite Fock space above can accordingly be recast as a *single* Fock space for "particles" only, with energies now covariantly running over the *whole* spectrum of positive and negative eigenvalues. Let the latter (strictly covariant) Fock space be denoted by $\mathscr{F}^\circ$. In line with this, the Stüeckelberg–Feynman approach can also enable one to think of a *covariant* charge conjugation which may *globally* interchange an either positive- or negative-energy fermion and an either positive- or negative-energy antifermion: in principle, this is *not* just the same as "particle" $\rightleftarrows$ "hole" conjugation (which is merely a noncovariant operation interchanging positive-energy fermions and antifermions). There is, however, one apparent difficulty standing in the way. As a matter of fact, the single Fock space $\mathscr{F}^\circ$ can *equally* pertain (covariantly) to either identical fermions or identical antifermions: according to the Stüeckelberg–Feynman views, a complete set of $\mathscr{F}^\circ$ kets (bras) for fermions amounts to a complete set of $\mathscr{F}^\circ$ bras (kets) for antifermions. Such a difficulty can actually be overcome by carefully reexamining the Stüeckelberg–Feynman approach in classical relativistic terms. Let $-p^\mu = m(-u^\mu)$ ($\mu = 0,1,2,3$; metric: $+ - - -$) be the four-momentum of a negative-energy particle of proper (i.e., covariant) mass m (> 0) and four-velocity $-u^{\mu} = -dx^{\mu}/ds$ ($-dx^0 < 0$). Since the equivalent positive-energy antiparticle, of four-momentum p^{μ}, is covering just the same world line in the opposite direction, $ds \rightarrow -ds$, the "slope" $-u^{\mu}$ of that world line cannot be affected by the reinterpretative procedure, $(-dx^{\mu})/ds = dx^{\mu}/(-ds)$. Strictly speaking, therefore, the procedure is *such that* $-p^{\mu} \rightarrow p^{\mu} \Rightarrow m \rightarrow -m$. On these grounds, one may state that a Dirac fermion and a Dirac antifermion can *covariantly* be distinguished by the (opposite) sign of their proper mass; and the covariant charge conjugation we are looking for is then to be identified with *proper-mass conjugation*.[13–16] As $\mathcal{F}^{\circ}$ can equally refer to either fermions or antifermions (with both positive and negative energies), we must expect it to be left invariant by proper-mass conjugation: this corresponds to the fact that the proper-mass sign in the Dirac equation is irrelevant. To get really a *non*trivial definition of a covariant charge conjugation, one should therefore *double* $\mathcal{F}^{\circ}$, by giving it some "label" that may specifically tell which of the two proper-mass signs is in turn being considered. For this purpose it is appropriate to introduce two (orthogonal) unit internal state vectors, $|f\rangle$ and $|\bar{f}\rangle$, as eigenvectors of a (covariant) one-particle proper-mass operator M with eigenvalues $+m$ and $-m$:

$$M|f\rangle = +m|f\rangle, \qquad M|\bar{f}\rangle = -m|\bar{f}\rangle \tag{1.1}$$

Let $\mathcal{S}_{\text{in}}$ be the two-dimensional internal space that is spanned by such eigenvectors. Then, a "dressed" generalized Fock space $\mathcal{F}$ can be built from the "bare" one $\mathcal{F}^{\circ}$, such that

$$\mathcal{F} \equiv \mathcal{F}^{\circ} \otimes \mathcal{S}_{\text{in}} \tag{1.2}$$

In this way the complete set of $\mathcal{F}^{\circ}$ kets (bras) may just undergo a *doubling* into a "Dirac fermionic" set, covariantly labeled by $|f\rangle$ ($\langle f|$), plus a "Dirac antifermionic" one, covariantly labeled by $|\bar{f}\rangle$ ($\langle \bar{f}|$) (with an energy range, in either case, still including both positive and negative eigenvalues); and the covariant charge conjugation may just be represented by a unitary and Hermitian operator C essentially acting in $\mathcal{S}_{\text{in}}$ and anticommuting with M:

$$C|f\rangle = |\bar{f}\rangle, \qquad C|\bar{f}\rangle = |f\rangle \qquad (C^{-1} = C^{\dagger} = C) \tag{1.3}$$

What such a doubling involves can be fully understood by coming back to the "particle–hole" language: it provides two *alternative* (equivalent) Dirac pictures where one is unambiguously choosing *either* "particle" = fermion and "hole" = antifermion, *or* "particle" = antifermion and "hole" = fermion, respectively. These are two *proper-mass-conjugated* descriptions, to be associated with two opposite-proper-mass Dirac free-field equations like

$$i\gamma^{\mu}\partial_{\mu}\psi_f = +m\psi_f, \qquad i\gamma^{\mu}\partial_{\mu}\,\psi_{\bar{f}} = -\,m\psi_{\bar{f}} \tag{1.4}$$

($\hbar = c = 1$; $\gamma^{0\dagger} = \gamma^0$, $\gamma^{k\dagger} = -\,\gamma^k$, $k = 1, 2, 3$) where $\psi_{\bar{f}}$ should consistently stand for the *proper-mass-conjugated* counterpart of ψ_f. Both field equations (1.4) are equally allowable within $\mathscr{F}^{\circ}$; and the dressed Fock space (1.2) should go along with a *double*-structured, dressed field operator of the type

$$\Psi(x) = \psi_f(x)\langle f| + \psi_{\bar{f}}(x)\langle \bar{f}| \tag{1.5}$$

($x \equiv x^{\mu}$). This is a Lorentz four-spinor, also looking like an $\mathscr{S}_{\text{in}}$ (bra) vector of "Dirac components" $\psi_f(x)$ and $\psi_{\bar{f}}(x)$ (whose orthogonality in $\mathscr{S}_{\text{in}}$ is just ensured by their being two proper-mass eigenfields with different eigenvalues). The field component $\psi_f(x)\langle f|$ can covariantly annihilate (either positive- or negative-energy) *Dirac fermions*, and the same holds for $\psi_{\bar{f}}(x)\langle \bar{f}|$ as regards (either positive- or negative-energy) *Dirac antifermions*. According to (1.3), the C-conjugate field operator reads

$$\Psi^{(C)}(x) \equiv \Psi(x)C = \psi_f(x)\langle \bar{f}| + \psi_{\bar{f}}(x)\langle f| \tag{1.6}$$

and a glance at (1.6) shows that applying C may equivalently be seen as putting

$$C\!: \qquad \psi_f(x) \rightleftarrows \psi_{\bar{f}}(x) \tag{1.7}$$

This just implies that $\psi_{\bar{f}}(x)$ is to be *covariantly* obtained from $\psi_f(x)$ (up to a phase factor) by applying proper-mass reversal to the Dirac equation:

$$\psi_{\bar{f}}(x) = \gamma^5\psi_f(x), \qquad \overline{\psi}_{\bar{f}}(x) = -\overline{\psi}_f(x)\gamma^5 \tag{1.8}$$

($\overline{\psi} = \psi^{\dagger}\gamma^0$; $\gamma^5 \equiv i\gamma^0\gamma^1\gamma^2\gamma^3$). In line with (1.8) (and with the fact that C is defined in $\mathscr{S}_{\text{in}}$) the standard Fourier expansions of $\psi_f(x)$ and $\psi_{\bar{f}}(x)$ must contain *identical* "particle"-annihilation operators, of the type $a(\mathbf{p}, \sigma)$, as well as *identical* "hole"-creation operators, of the type $a^{\text{h}\dagger}(\mathbf{p}, \sigma)$ (σ being the helicity variable). This is admissible, because $\psi_f(x)$ and $\psi_{\bar{f}}(x)$ belong to *alternative* (proper-mass-conjugated) Dirac pictures where one has either "particle" = fermion and "hole" = antifermion, or "particle" = antifermion and "hole" = fermion, respectively, each single picture being already able (independently of the other) to account for the creation or annihilation of a "particle"–"hole" pair. Here $\psi_{\bar{f}}(x)$ has nothing to do with the customary "hole" field that is (noncovariantly) obtained by applying "particle" $\rightleftarrows$ "hole" conjugation (i.e., $a \to a^{\text{h}}$, $a^{\text{h}\dagger} \to a^{\dagger}$): the latter field can still be encountered, too, but only *within* either picture, as a mere result of normal ordering. By exploiting Eqs. (1.8) and introducing the adjoint field operator

$$\overline{\Psi}(x) = |f\rangle\overline{\psi}_f(x) + |\bar{f}\rangle\overline{\psi}_{\bar{f}}(x) \tag{1.9}$$

one can compactly write

$$\Psi^{(C)}(x) = \gamma^5\Psi(x), \qquad \overline{\Psi}^{(C)}(x) = -\,\overline{\Psi}(x)\gamma^5 \tag{1.10}$$

This strictly defines the effective action of chirality γ^5 as covariant charge conjugation.

Another remarkable $\mathcal{S}_{\text{in}}$ basis can be obtained from $(|f\rangle, |\bar{f}\rangle)$ by performing the rotation

$$\begin{aligned} |f\rangle &= 2^{-1/2}(|f^{\text{ch}}\rangle + |\bar{f}^{\text{ch}}\rangle) \\ |\bar{f}\rangle &= 2^{-1/2}(-|f^{\text{ch}}\rangle + |\bar{f}^{\text{ch}}\rangle) \end{aligned} \tag{1.11}$$

The operator C is made diagonal in such a basis:

$$C|f^{\text{ch}}\rangle = -|f^{\text{ch}}\rangle, \qquad C|\bar{f}^{\text{ch}}\rangle = |\bar{f}^{\text{ch}}\rangle \tag{1.12}$$

Thus a further (unitary and Hermitian) operator, say P_{in}, can be introduced in $\mathcal{S}_{\text{in}}$, having the property of interchanging $|f^{\text{ch}}\rangle$ and $|\bar{f}^{\text{ch}}\rangle$,

$$P_{\text{in}}|f^{\text{ch}}\rangle = |\bar{f}^{\text{ch}}\rangle, \qquad P_{\text{in}}|\bar{f}^{\text{ch}}\rangle = |f^{\text{ch}}\rangle \qquad (P_{\text{in}}^{-1} = P_{\text{in}}^{\dagger} = P_{\text{in}}) \tag{1.13}$$

and, in its turn, being diagonal in the basis $(|f\rangle, |\bar{f}\rangle)$. More specifically, since

$$P_{\text{in}}|f\rangle = |f\rangle, \qquad P_{\text{in}}|\bar{f}\rangle = -\,|\bar{f}\rangle \tag{1.14}$$

one has that P_{in} may just be interpreted (apart from a phase constant $\eta = \pm 1$) as standing for an "intrinsic parity" covariant operator: it should be identified with that factor of the parity operator P ($\equiv P_{\text{ex}}P_{\text{in}} = P_{\text{in}}P_{\text{ex}}$) which properly acts in $\mathcal{S}_{\text{in}}$ (the other factor, P_{ex}, properly acting in $\mathcal{F}^{\circ}$). In the new $\mathcal{S}_{\text{in}}$ basis, the field $\Psi(x)$ reads

$$\Psi(x) = \chi_f(x)\langle f^{\text{ch}}| + \chi_{\bar{f}}(x)\langle \bar{f}^{\text{ch}}| \tag{1.15}$$

where

$$\chi_f(x) \equiv 2^{-1/2}(1 - \gamma^5)\psi_f(x), \qquad \chi_{\bar{f}}(x) \equiv 2^{-1/2}(1 + \gamma^5)\psi_{\bar{f}}(x) \tag{1.16}$$

and

$$\psi_f = 2^{-1/2}(\chi_f + \chi_{\bar{f}}), \qquad \psi_{\bar{f}} = 2^{-1/2}(-\chi_f + \chi_{\bar{f}}) \tag{1.17}$$

In this way one can naturally obtain two (massive) *chiral* fields χ_f and $\chi_{\bar{f}}$ with *opposite* chiralities, on the same footing as the two Dirac fields ψ_f and $\psi_{\bar{f}}$. Such an outcome *a fortiori* makes sense in the zero-mass limit, where it is made both *automatic* and *universal* to get only *two* (rather than four) independent chiral-field solutions being covariantly available in all: these are a *left-handed* one, just for a (massless) fermion, and a *right-handed* one, just for a (massless) antifermion. Hence, in particular, it can strictly be stated (without the help of any ad hoc prescription) that *if the neutrino proper mass is assumed to be vanishing, then no right-handed neutrinos and no left-*

handed antineutrinos may, in principle, be expected to exist (in addition to left-handed neutrinos and right-handed antineutrinos).[18] This is also to be related to the more general fact that, by virtue of (1.17), the "maximally *P*-violating" Dirac-field *V–A* weak current may now be given a *natural* place at the theoretic level, provided it is *rigorously* recast as a *chiral-field pure vector current* according to the formula.[2]

$$\overline{\psi}_f^{(a)}\gamma^\mu(1-\gamma^5)\psi_f^{(b)} \equiv \overline{\chi}_f^{(a)}\gamma^\mu\chi_f^{(b)} \tag{1.18}$$

($\overline{\chi} = \chi^\dagger\gamma^0$) (the superscripts a, b covariantly labeling the fermions connected). The same could alternatively be said for a $V + A$ current connecting two antifermions a, b:

$$\overline{\psi}_{\bar f}^{(a)}\gamma^\mu(1+\gamma^5)\psi_{\bar f}^{(b)} \equiv \overline{\chi}_{\bar f}^{(a)}\gamma^\mu\chi_{\bar f}^{(b)} \tag{1.19}$$

So, *P* symmetry itself (with *P* still represented by γ^0) can be formally restored via either (1.18) or (1.19) (since γ^0 is now to be applied *directly* on chiral fields, rather than on Dirac fields); and the new currents are further left *unvaried* by the covariant charge conjugation *C* (since both χ_f and $\chi_{\bar f}$ are *C*-eigenfields). As discussed in ref. 1, such recovered *P* symmetry (on passing to the generalized formalism in hand) actually corresponds to the phenomenological "*CP* symmetry" (according to the standard formalism): here the left–right asymmetry of weak fermionic processes amounts to a P_{ex} violation (P_{ex} being the mere "external" parity acting in $\mathscr{F}^\circ$) and not just to a P ($= P_{in}P_{ex}$) violation; likewise the standard "charge-conjugation failure" characterizing those processes amounts to a P_{in} failure (such that the product $P_{in}P_{ex}$ is still a symmetry operation).

A fuller insight into this can be gained by introducing two one-particle "charge" operators Q and Q^{ch}, the former being diagonal (with opposite nonzero eigenvalues) in the "Dirac" $\mathscr{S}_{in}$ basis ($|f\rangle$, $|\bar f\rangle$) and the latter being the same in the "chiral" $\mathscr{S}_{in}$ basis ($|f^{ch}\rangle$, $|\bar f^{ch}\rangle$): they are such that

$$CQ = -QC, \qquad P_{in}Q = QP_{in} \tag{1.20}$$

and

$$P_{in}Q^{ch} = -Q^{ch}P_{in}, \qquad CQ^{ch} = Q^{ch}C \tag{1.21}$$

Thus Q behaves like a *scalar* charge (reversed by C) and Q^{ch} like a *pseudoscalar* charge (reversed by P_{in}); and one has that C and P_{in} properly stand for *scalar*- and *pseudoscalar-charge conjugation* operators, respectively. Moreover, it follows from Eqs. (1.3) and (1.13) that the internal states ($|f\rangle$, $|\bar f\rangle$) look like *pure* scalar-charge conjugated eigenstates, whereas the internal

[2] See ref. 19. for the well-established *universality* of the pure chiral field formal structure in the weak phenomenology of leptons and quarks.

states ($|f^{ch}\rangle$, $|\bar{f}^{ch}\rangle$) look like *pure* pseudoscalar-charge conjugated eigenstates. This is directly connected with the fact that Q and Q^{ch} are *anticommuting* operators,

$$QQ^{ch} + Q^{ch}Q = 0 \tag{1.22}$$

though their squares clearly satisfy the commutation relations

$$[Q^2, Q^{ch}] = [Q^{ch2}, Q] = 0 \tag{1.23}$$

Each of the two charges Q and Q^{ch}, if singly applied (from the right) to the field $\Psi(x)$, is able to *superselect* that internal representation of $\Psi(x)$—either (1.5) or (1.15)—diagonalizing it. So, the same fermion–antifermion pair may in turn be suitably described by the Dirac or the chiral $\mathscr{S}_{in}$ basis according to whether a charge Q or Q^{ch} is alternately involved.[20] Hence it can be argued that the "true" fermion $\rightarrow$ antifermion covariant conjugation should generally be identified with CP_{in}, even though CP_{in} is just reducible to C when acting on $|f\rangle$ and to P_{in} when acting on $|f^{ch}\rangle$:

$$CP_{in}|f\rangle = C|f\rangle, \qquad CP_{in}|f^{ch}\rangle = P_{in}|f^{ch}\rangle \tag{1.24}$$

($CP_{in} = -P_{in}C$). In the former case the fermion behaves like a scalar-charge object, in the latter like a pseudoscalar-charge object; but in *both* individual cases (and not only in the former one) P mirror symmetry may be strictly respected:[21] as particularly regards the latter case, the state of a fermion at rest appears to be *no longer* a P-eigenstate, and P plays also an *internal* role as "(pseudoscalar)-charge conjugation" (in place of C).

The field $\Psi(x)$, as given by either (1.5) or (1.15), obeys the generalized free Dirac equation

$$i\gamma^{\mu}\, \partial_{\mu}\Psi(x) = \Psi(x)M \tag{1.25}$$

M being the one-particle proper-mass operator defined by (1.1). A comparison of (1.1) with (1.14) makes it possible to recast this equation in the more convenient form

$$i\gamma^{\mu}\, \partial_{\mu}\Psi(x) = |m|\Psi^{(P_{in})}(x) \tag{1.26}$$

where

$$\Psi^{(P_{in})}(x) \equiv \Psi(x)P_{in} = \psi_f(x)\langle f| - \psi_{\bar{f}}(x)\langle \bar{f}| \tag{1.27}$$

The field equation (1.26) (as well as its P_{in}-conjugated counterpart) is derivable from the real free Lagrangian density

$$\begin{aligned}\mathscr{L}(\Psi, \Psi^{(P_{in})}, \overline{\Psi}, \overline{\Psi}^{(P_{in})}, \ldots; |m|) = \tfrac{1}{4}[i(\overline{\Psi}\gamma^{\mu}\partial_{\mu}\Psi + \overline{\Psi}^{(P_{in})}\gamma^{\mu}\partial_{\mu}\Psi^{(P_{in})}) + \text{H.c.}]\\ - \tfrac{1}{2}|m|(\overline{\Psi}\Psi^{(P_{in})} + \overline{\Psi}^{(P_{in})}\Psi)\end{aligned} \tag{1.28}$$

where $\overline{\Psi}^{(P_{\rm in})} = P_{\rm in}\overline{\Psi}$. A glance at (1.28) shows that $\mathscr{L}$ is both manifestly $P_{\rm in}$-invariant and P-invariant, with P acting as usual:

$$P: \quad \partial_\mu \to \partial^\mu, \quad \gamma^\mu \to \gamma^0\gamma^\mu\gamma^0 \tag{1.29}$$

This is a quite covariant result, *without regard* to the special $\mathscr{S}_{\rm in}$-representation being utilized for the fields $\Psi, \Psi^{(P_{\rm in})}$, $\overline{\Psi}$, and $\overline{\Psi}^{(P_{\rm in})}$; so that parity invariance consistently holds even in the *chiral* $\mathscr{S}_{\rm in}$-representation. Global phase invariance of $\mathscr{L}$ yields a manifestly $P_{\rm in}$-invariant, conserved free current like

$$J \equiv J^\mu = \tfrac{1}{2}[\overline{\Psi}\gamma^\mu\Psi + \overline{\Psi}^{(P_{\rm in})}\gamma^\mu\Psi^{(P_{\rm in})}] \tag{1.30}$$

If the closure relation $|f\rangle\langle f| + |\bar{f}\rangle\langle\bar{f}| = 1$ is exploited (throughout this paper, the identity operator in $\mathscr{S}_{\rm in}$ will be simply denoted by 1), such a current can be reduced to the form

$$J^\mu = \overline{\psi}_f\gamma^\mu\psi_f = \overline{\psi}_{\bar{f}}\gamma^\mu\psi_{\bar{f}} = \tfrac{1}{2}[\overline{\chi}_f\gamma^\mu\chi_f + \overline{\chi}_{\bar{f}}\gamma^\mu\chi_{\bar{f}}] \tag{1.31}$$

acting merely in the bare Fock space $\mathscr{F}^\circ$. This form can be suitably "dressed" to give the two distinct, scalar- and pseudoscalar-charge, conserved free currents

$$\mathscr{J}^{(Q)} = QJ = JQ, \quad \mathscr{J}^{(Q^{\rm ch})} = Q^{\rm ch}J = JQ^{\rm ch} \tag{1.32}$$

being, according to (1.20) and (1.21), such that

$$C\mathscr{J}^{(Q)} = -\mathscr{J}^{(Q)}C, \qquad P_{\rm in}\mathscr{J}^{(Q^{\rm ch})} = -\mathscr{J}^{(Q^{\rm ch})}P_{\rm in} \tag{1.33}$$

and

$$P_{\rm in}\mathscr{J}^{(Q)} = \mathscr{J}^{(Q)}P_{\rm in}, \qquad C\mathscr{J}^{(Q^{\rm ch})} = \mathscr{J}^{(Q^{\rm ch})}C \tag{1.34}$$

Note, in particular, that the *vector* current $\mathscr{J}^{(Q)}$ commutes with $P_{\rm in}$, whereas the *axial-vector* current $\mathscr{J}^{(Q^{\rm ch})}$ anticommutes with $P_{\rm in}$. Such an outcome is covariant; so it may be extended to *all* vectors and axial vectors acted upon by $P_{\rm in}$. As for the scalar-charge current $\mathscr{J}^{(Q)}$, let q $(= \pm|q| \neq 0)$ and $-q$ be the eigenvalues of its "dressing" one-particle charge operator Q. Then Q can be expressed as

$$Q(|q|) = \pm|q|(\mathscr{P}_f - \mathscr{P}_{\bar{f}}) \tag{1.35}$$

with

$$\mathscr{P}_f \equiv |f\rangle\langle f|, \qquad \mathscr{P}_{\bar{f}} \equiv |\bar{f}\rangle\langle\bar{f}| \tag{1.36}$$

and $\mathscr{J}^{(Q)}$ can be given the *double*-structured, explicit form

$$\mathscr{J}^{(Q)} \equiv \mathscr{J}^{(Q)\mu} = q(\overline{\psi}_f\gamma^\mu\psi_f)\mathscr{P}_f + (-q)(\overline{\psi}_{\bar{f}}\gamma^\mu\psi_{\bar{f}})\mathscr{P}_{\bar{f}} \tag{1.37}$$

In line with the peculiar meaning of the dressed Fock space (1.2), this form joins two *alternative* (and equivalent) covariant currents—a "Dirac fermionic" one (marked by $\mathcal{P}_f$) and a "Dirac antifermionic" one (marked by $\mathcal{P}_{\bar{f}}$)—just pertaining to the two above equally allowable proper-mass-conjugated descriptions. Of course, *both* alternative pictures embodied in (1.37) are consistent with QED, since a Dirac current term like $\overline{\psi}\gamma^\mu\psi$ is left invariant by proper-mass reversal $\psi \to \gamma^5\psi$, $\overline{\psi} \to -\overline{\psi}\gamma^5$. On applying normal ordering, the "Dirac fermionic" sector of (1.37) can be made equivalent to a complete (antisymmetrized) "particle + hole" current where "particle" = fermion and "hole" = antifermion, the converse being true for the "Dirac antifermionic" sector; and within either of these normally ordered sectors (each one being marked by a *single* proper-mass sign) the standard Dirac charge conjugation can make its appearance again, with the noncovariant meaning of "particle" $\rightleftarrows$ "hole" conjugation. The Lagrangian density (1.28) may acquire also local invariance with respect to U(1) transformations generated by Q, provided it is supplemented by a minimal covariant coupling term like $-\mathcal{J}^{(Q)}A$, where $A \equiv A_\mu$ is a (massless) vector gauge field. This is a *double*-structured term as well: it actually joins two equivalent, and only alternative, coupling terms pertaining to the two ("Dirac fermionic" and "Dirac antifermionic") currents embodied in $\mathcal{J}^{(Q)}$. The term $-\mathcal{J}^{(Q)}A$ is left invariant by the chirality transformation $\psi \to \gamma^5\psi$, $\overline{\psi} \to -\overline{\psi}\gamma^5$; hence, the covariant scalar-charge conjugation C as just represented by γ^5 is always defined as a *symmetry* operation, which acts, namely, on the *whole* interacting system (A_μ included):

$$C\mathcal{J}^{(Q)}AC^\dagger = (-\mathcal{J}^{(Q)})(-A) \tag{1.38}$$

2. A NATURAL TWO-COMPONENT THEORY FOR A MASSLESS FERMION, WITH *P* MIRROR SYMMETRY RECOVERED

The "intrinsic parity" covariant operator defined by either (1.13) or (1.14) is such that (in the chiral $\mathcal{S}_{\text{in}}$ representation) one has

$$\Psi(x)P_{\text{in}} = \chi_f(x)\langle \bar{f}^{\text{ch}}| + \chi_{\bar{f}}(x)\langle f^{\text{ch}}| \tag{2.1}$$

A glance at (2.1) (and its adjoint) shows that P_{in} may equivalently be seen as yielding *chirality conjugation*:

$$P_{\text{in}}: \quad \chi_f(x) \rightleftarrows \chi_{\bar{f}}(x), \quad \overline{\chi}_f(x) \rightleftarrows \overline{\chi}_{\bar{f}}(x) \tag{2.2}$$

That plays a fundamental role in the zero-mass limiting case, where chirality becomes coincident with helicity (except for a change of sign from positive to negative energies) and the P_{in} operation then takes the special meaning of *helicity conjugation*.

On setting $|m| = 0$, the Lagrangian density (1.28) is reduced to

$$\mathcal{L}(|m| = 0) = \tfrac{1}{4}[i(\overline{\Psi}\gamma^{\mu}\partial_{\mu}\Psi + \overline{\Psi}^{(P_{\rm in})}\gamma^{\mu}\partial_{\mu}\Psi^{(P_{\rm in})}) + \text{H.c.}] \tag{2.3}$$

This form still has manifest covariance (extended to $\mathcal{S}_{\rm in}$) and clearly *retains* P invariance. The *twofold*—either "Dirac" or "chiral"—internal nature of a massive spin-1/2 fermion (antifermion) is now lost, since a massless (and definite-helicity) fermion (antifermion) can *only* exist as a definite-chirality particle: in the zero-mass case, the above two C-conjugated Dirac fields ψ_f and $\psi_{\bar{f}}$ can be involved *only as mixtures*—according to (1.17)—*of the single, definite-chirality (true fermionic and antifermionic) fields* χ_f and $\chi_{\bar{f}}$. So, if we allow *ab initio* for the chiral condition (1.16) and substitute it into (2.3), we may exploit the closure relation $|f^{\rm ch}\rangle\langle f^{\rm ch}| + |\bar{f}^{\rm ch}\rangle\langle\bar{f}^{\rm ch}| = 1$ to recast (2.3) in the following "undressed" form[18]:

$$\mathcal{L}(|m| = 0) = \tfrac{1}{4}[i\,(\overline{\chi}_f\gamma^{\mu}\partial_{\mu}\chi_f + \overline{\chi}_{\bar{f}}\gamma^{\mu}\partial_{\mu}\chi_{\bar{f}}) + \text{H.c.}] \tag{2.4}$$

where χ_f and $\chi_{\bar{f}}$ are now two quite distinct field operators that annihilate *opposite*-helicity one-particle states (with either positive or negative energies) belonging to the pure covariant Fock space $\mathcal{F}^{\circ}(|m| = 0)$. The Lagrangian density (2.4) possesses manifest invariance under (2.2). Moreover, the *retrieval* of P invariance in (2.4) can directly be checked by considering that applying (1.29) gives

$$P\!: \quad \overline{\chi}\gamma^{\mu}\partial_{\mu}\chi \to \overline{\chi}\gamma^{0}\gamma^{\mu}\partial^{\mu}\gamma^{0}\chi = \overline{\chi}\gamma^{\mu}\partial_{\mu}\chi \tag{2.5}$$

One essential physical aspect of the recovered P symmetry lies in the fact that, due to the former equation in (1.8), we now have

$$\chi_{\bar{f}} = \xi_f \equiv 2^{-1/2}(1 + \gamma^5)\psi_f, \qquad \chi_f = -\,\xi_{\bar{f}} \equiv 2^{-1/2}(1 - \gamma^5)\psi_{\bar{f}} \tag{2.6}$$

ξ_f and $\xi_{\bar{f}}$ are the "missing" helicity-conjugated counterparts of χ_f and $\chi_{\bar{f}}$ in the ordinary approach. The effective operation interchanging a massless fermion and antifermion may thus take the more significant form

$$P_{\rm in}\!: \quad \chi_f(x) = -\xi_{\bar{f}}(x) \rightleftarrows \chi_{\bar{f}}(x) = \xi_f(x) \tag{2.7}$$

where $P_{\rm in}$ manifestly acts as a net helicity-conjugation operator (and $P_{\rm in}$ symmetry then becomes manifestly equivalent to helicity-conjugation symmetry). So, by using Weyl's γ-matrix representation, one naturally obtains a *two*-component fermion antifermion field scheme where the helicity-conjugated counterpart of the actual fermion (antifermion) field is not "missing" at all, but *already coincides* with the actual antifermion (fermion) field itself! Such an outcome can be summed up by the following statement: *applying an effective particle ⇄ antiparticle conjugation to either a left-handed massless fermion or a right-handed massless antifermion merely means applying helicity conjugation to it, without any real failure of P mirror symmetry.* This is to be supplemented by the fact that we may also speak of a "recovered C

symmetry" (in a way reminiscent of the Majorana neutrino[22]) provided C stands for the *covariant* scalar-charge conjugation represented by γ^5. All that, on the other hand, is not an *ad hoc* theory for the only neutrino (such as the standard two-component scheme[8–11]), but universally applies to *any spin*-1/2 *point fermion in its massless original version*.

Global invariance of (2.4) under the chiral group $U^{(\gamma^5)}(1)$ yields a conserved free chiral current like

$$J^{\mathrm{ch}(\gamma^5)} \equiv \tfrac{1}{2}[\bar{\chi}_f\gamma^\mu q^{\mathrm{ch}}(-\gamma^5)\chi_f + \bar{\chi}_{\bar f}\gamma^\mu q^{\mathrm{ch}}(-\gamma^5)\chi_{\bar f}]$$
$$= \tfrac{1}{2}q^{\mathrm{ch}}(\bar{\chi}_f\gamma^\mu\chi_f - \bar{\chi}_{\bar f}\gamma^\mu\chi_{\bar f}) \tag{2.8}$$

Such a current carries the *pseudoscalar* undressed charge (operator) $q^{\mathrm{ch}}(-\gamma^5)$, of which the two fields χ_f and $\chi_{\bar f}$ are conjugated eigenfields (with opposite eigenvalues, q^{ch} and $-q^{\mathrm{ch}}$). Apart from the explicit presence of a charge operator like $q^{\mathrm{ch}}(-\gamma^5)$, a current of the form (2.8) is strictly bound to be a pseudoscalar-charge current by the fact that

$$P_{\mathrm{in}}:\quad \bar{\chi}_f\gamma^\mu\chi_f \rightleftarrows \bar{\chi}_{\bar f}\gamma^\mu\chi_{\bar f} \tag{2.9}$$

This *prevents* χ_f and $\chi_{\bar f}$ from being also *scalar-charge-conjugated eigenfields* (with opposite eigenvalues), as the related current would be again of the type (2.8) and would again be inverted by P_{in}. Yet we may think of a further conserved free current like

$$J^{\mathrm{ch}} \equiv \pm|q|\tfrac{1}{2}(\bar{\chi}_f\gamma^\mu\chi_f + \bar{\chi}_{\bar f}\gamma^\mu\chi_{\bar f}) \tag{2.10}$$

where the carried c-number $\pm\,|q|$ ($\neq 0$) is the *same* for χ_f *as* for $\chi_{\bar f}$. Such a current stems from global invariance under the group U(1) generated by either $+|q|$ or $-|q|$, and the value $\pm|q|$ appearing in (2.10) can but be interpreted as a conventional root of $|q|^2$. We can thus argue that the retrieval of parity symmetry for a left-handed massless fermion and a right-handed massless antifermion should be universally due to *their being two pseudoscalar-charge-conjugated eigenstates* (which are turned into each other by ordinary space inversion): if also some scalar (additive) charge is assumed to belong to them, it can be at most determined *only in its* (squared) *magnitude* $|q|^2$, whose (conventionally chosen) root—either $+|q|$ or $-|q|$ for *them both*—essentially gives the mere *absolute strength* of that charge (in line with the general *anticommutivity* property of scalar- and pseudoscalar-charge conjugations C and P_{in}).

In an overall view (with and without mass) such a pure *theoretical* understanding of the "maximal parity-violation" effect makes it appropriate to attempt some further insight into the "standard model," whose minimal formulation,[23–25] as is well known, can but (phenomenologically) *postulate* the "parity-violating" nature of the weak-isospin fermionic current, leaving quite unsolved the question of its origin.[4] These new prospects are opened by the present paper, which deals specifically with the leptonic sector of the electroweak scheme.

3. A LINK BETWEEN THE APPEARANCE OF FERMION MASSES AND THE INTERNAL PRESENCE OF SCALAR CHARGES GENERATING LOCAL GAUGE SYMMETRIES

The *twofold*—either Dirac or chiral—massive-fermion model sketched in Section 1 shows one peculiar feature that is worth further analyzing. Due to the *anticommutivity* property (1.22), which holds within the generalized Fock space (1.2), the presence of an effective *superselection rule* is now made a *non*trivial requirement for a given charge to be enabled to generate a local gauge symmetry: for instance, as already pointed out, a *superselecting* mechanism is switched on for a scalar charge as soon as the fermion–antifermion generalized field $\Psi(x)$ is applied to (from the right) by a suitable (one–particle) charge *operator* of the Q type (defined in $\mathscr{S}_{\text{in}}$) whose action just selects the *Dirac* $\mathscr{S}_{\text{in}}$ representation of $\Psi(x)$ (i.e., the one diagonalizing the scalar charge). The extreme consequences of this feature can be found on passing to the zero-mass case; they are concerned with *scalar*-charge local gauge symmetries and particularly refer to the electromagnetic local gauge symmetry. The following statement can specifically be proved: *An electric charge being carried by a massless fermion (antifermion) should be unable yet to generate a local gauge coupling; it can actually do so only if the fermion (antifermion) is made massive*. The proof directly stems from the above P-symmetric model of a massless fermion–antifermion pair. According to such a model, if an electrically charged spin-1/2 fermion and the associated antifermion really happened to be massless, they should be described by two (left- and right-handed) chiral fields like χ_f and $\chi_{\bar{f}}$, respectively. These, however, cannot stand for two *scalar-charge-conjugated eigenfields*, since they are rather the net *helicity-conjugated* counterparts of each other. The fermion and antifermion in question would therefore behave as carrying at most an electric charge *definite but in magnitude* (the same for them both). Hence, *no* superselection rule could ever apply to that charge—whose eigenfields would be only *mixtures* of χ_f and $\chi_{\bar{f}}$ according to (1.17)—and *no* mechanism could ever *diagonalize* it (thus enabling it to generate the electromagnetic gauge coupling).

This puts one in a position to set *anew* the question about spontaneous symmetry breaking (SSB) and fermion masses. In the standard approach to the electroweak model,[24,25] the final charged leptons and quarks look massive but for an external "accident," connected with their hypothetical couplings to the Higgs boson[26,27]: in the absence of such couplings (which are conjectured ad hoc to accommodate theory to experience) they might well have been left massless with no matter of principle. In the present approach, on the contrary, the mass generation for charged leptons and quarks does become an *essential intrinsic requirement* of the model, as just follows from the theorem above. Herein, the appearance of a mass for an electrically charged fermion looks rather like an *unavoidable* outcome of the *emerging* dynamical nature itself of the electric charge carried by that fermion. Hence a more appropriate general mechanism of SSB should be expected to work, somehow *free* from the constraint of a strict "external" real agent responsible for it.

4. WEAK ISOSPIN REDEFINED

In the standard formulation, the one-particle weak-isospin operator $\mathbf{T}^{\mathrm{W}}$ is ad hoc imposed to act merely on left-handed, and not on right-handed, massless fermion fields. That appears to be a theoretically unnatural prescription, which is clearly needed, however, to allow for the "maximal *P*-violation" effect. As a final outcome after the mass appearance, those which (according to standard views) are merely the two chiral projections of one and the *same* (massive) Dirac field are caused to fall within *different* weak-isospin representations.

The reason for such an *accommodation* actually drops on passing to the strictly covariant fermion–antifermion quantum field framework in hand: it naturally includes a massless scheme where just *two* (rather than four) independent chiral-field solutions turn out to be available in all, namely a right-handed one for the *only* antifermion and a left-handed one for the *only* fermion.

In line with what has been argued in Section 2, each individual component of $\mathbf{T}^{\mathrm{W}}$ is now required to behave as a *pseudoscalar* quantity in ordinary space, so that $\mathbf{T}^{\mathrm{W}}$ itself should be inverted by pseudoscalar-charge conjugation:

$$P_{\mathrm{in}}: \quad \mathbf{T}^{\mathrm{W}} \to -\mathbf{T}^{\mathrm{W}} \tag{4.1}$$

Let then $\mathbf{T}^{\mathrm{W}}$ be conveniently redefined as

$$\mathbf{T}^{\mathrm{W}} = \tfrac{1}{2}\boldsymbol{\tau}(-\gamma^5) \tag{4.2}$$

$\frac{1}{2}\boldsymbol{\tau}$ looks like an ordinary isospin-1/2 operator (and the $\boldsymbol{\tau}$ components τ_1, τ_2, and τ_3 are formally identical with the three Pauli matrices). So, due to the presence of $(-\gamma^5)$, the new weak-isospin operator (4.2) cannot generate any SU(2) transformation group: from it, rather, *two* separate (left- and right-handed) SU(2) generators can proceed as soon as $(-\gamma^5)$ is diagonalized and its (either positive or negative) eigenvalue is specified. More precisely, let us put

$$\mathbf{T}^{\mathrm{W}} = \mathbf{T}_L + \mathbf{T}_R \tag{4.3}$$

where

$$\mathbf{T}_L = \tfrac{1}{2}\boldsymbol{\tau}(-\gamma^5)X_L, \qquad \mathbf{T}_R = \tfrac{1}{2}\boldsymbol{\tau}(-\gamma^5)X_R \tag{4.4}$$
$$X_L \equiv \tfrac{1}{2}(1 - \gamma^5), \qquad X_R \equiv \tfrac{1}{2}(1 + \gamma^5)$$

and

$$P_{\mathrm{in}}: \quad \mathbf{T}_L \to -\mathbf{T}_R, \quad \mathbf{T}_R \to -\mathbf{T}_L \tag{4.5}$$

It is immediate then to realize that the two operators $\mathbf{T}_L$ and $\mathbf{T}_R$, if taken individually, are just able to generate two distinct SU(2) groups, $SU(2)_L$ and $SU(2)_R$, whose fundamental representations (the former acting on left-handed chiral fields and the latter on right-handed ones) are marked by the *effective* isospin generators $\frac{1}{2}\boldsymbol{\tau}$ and $\frac{1}{2}(-\boldsymbol{\tau})$, respectively. One may accordingly think also of two distinct *effective* isospin spaces, associated with $\frac{1}{2}\boldsymbol{\tau}$ and $\frac{1}{2}(-\boldsymbol{\tau})$, and build *either* a neutrino–electron (massless) $SU(2)_L$ isodoublet

$$D_L = \begin{pmatrix} \chi_\nu \\ \chi_e \end{pmatrix} \tag{4.6}$$

or an antineutrino–positron (massless) $SU(2)_R$ isodoublet

$$D_R = \begin{pmatrix} \chi_{\bar\nu} \\ \chi_{\bar e} \end{pmatrix} \tag{4.7}$$

where

$$T_{3_L}\begin{pmatrix} \chi_\nu \\ \chi_e \end{pmatrix} = \frac{1}{2}\,\tau_3\begin{pmatrix} \chi_\nu \\ \chi_e \end{pmatrix}, \qquad T_{3R}\begin{pmatrix} \chi_{\bar\nu} \\ \chi_{\bar e} \end{pmatrix} = \frac{1}{2}\,(-\tau_3)\begin{pmatrix} \chi_{\bar\nu} \\ \chi_{\bar e} \end{pmatrix} \tag{4.8}$$

(note that the definitions of D_L and D_R may look symmetrical, just because they do *not* refer to the same effective isospin group). These turn out to be two alternatively allowable (left- and right-handed) covariant pictures. In the former picture, it is just $\chi_{\bar\nu}$ and $\chi_{\bar e}$ that enter as $SU(2)_L$ isosinglets (in place of the ordinary, neutrino and electron, right-handed field solutions ξ_ν and ξ_e); and in the latter picture, similarly, it is just χ_ν and χ_e that enter as $SU(2)_R$ isosinglets:

$$\mathbf{T}_L D_R = 0, \qquad \mathbf{T}_R D_L = 0 \tag{4.9}$$

In view of all this, a *double*—either $SU(2)_L$ or $SU(2)_R$—variety of weak-isospin SU(2) representations may be conceived of. This is indeed a general *intrinsic* feature of the redefined weak isospin, which directly proceeds from (4.2) and holds regardless of the zero-mass condition.

5. A PARITY-INVARIANT REFORMULATION OF THE STANDARD LEPTONIC GAUGE SCHEME IN THE ABSENCE OF MASSES

We may begin reformulating the leptonic sector of the electroweak theory by setting, as usual, all masses equal to zero. It should be emphasized, however, that here, unlike in the ordinary version,[23–25] such a condition *prevents* charged leptons from already being true eigenstates of electric charge.

In dealing with the strict massless case, we can now rely upon individual lepton–antilepton free Lagrangian densities of the type (2.4). If we confine ourselves to the electron (antielectron) and the related neutrino (antineutrino), we may thus start off by putting

$$\mathscr{L}_{\nu\bar{\nu}} = \tfrac{1}{4}[i(\bar{\chi}_\nu\gamma^\mu\partial_\mu\chi_\nu + \bar{\chi}_{\bar{\nu}}\gamma^\mu\partial_\mu\chi_{\bar{\nu}}) + \text{H.c.}] \tag{5.1}$$

as well as

$$\mathscr{L}_{e\bar{e}} = \tfrac{1}{4}[i(\bar{\chi}_e\gamma^\mu\partial_\mu\chi_e + \bar{\chi}_{\bar{e}}\gamma^\mu\partial_\mu\chi_{\bar{e}}) + \text{H.c.}] \tag{5.2}$$

The Lagrangian density (5.1) naturally yields the neutrino two-component scheme without involving at all any failure of P mirror symmetry. A similar property holds for the electron–antielectron Lagrangian density (5.2), whose difference from (5.1) lies merely in the fact that χ_e and $\chi_{\bar{e}}$ are two electrically charged chiral fields. These, however, cannot be taken as just being two (conjugated) *eigenfields* of electric charge. What can at most be said is that *as long as the electron and positron are strictly assumed to be massless, the (chiral) fields associated with them should be ascribed an electric charge of a definite squared magnitude,* $|e|^2$, *but of an as yet quite indefinite sign*. The root of $|e|^2$—which may conventionally be chosen to be either $-|e|$ or $+|e|$—gives essentially the *absolute strength* of that charge; and so *the same* $|e|^2$ *root should together be assignable to both fields* χ_e *and* $\chi_{\bar{e}}$. This corresponds to the existence of a massless conserved free current like (2.10) of the general type

$$J^{\text{ch}}(\mp|e|) \equiv \mp|e|\tfrac{1}{2}(\bar{\chi}_e\gamma^\mu\chi_e + \bar{\chi}_{\bar{e}}\gamma^\mu\chi_{\bar{e}}) \tag{5.3}$$

where the minus or plus sign holds according to the choice made.

Let us take the whole free Lagrangian density

$$\mathscr{L} = \mathscr{L}_{\nu\bar{\nu}} + \mathscr{L}_{e\bar{e}} \tag{5.4}$$

which, like (5.1) and (5.2), is both P- and P_{in}-invariant, with P and P_{in} defined by (2.5) and (2.2). If we choose the minus sign in (5.3)—so as to make the $|e|^2$ root *numerically* coincident with the electric charge $-e$ of the actual (massive) electron—it becomes appropriate to join together the two fields χ_ν and χ_e into a weak isospin *doublet* D_L as given by (4.6). Doing like this, we are automatically selecting the *left-handed*, i.e., $SU(2)_L$, variety of the weak isospin representations. Within it, on the other hand, both (positron and antineutrino) fields $\chi_{\bar{e}}$ and $\chi_{\bar{\nu}}$ are to be classified as $SU(2)_L$ *singlets*, just corresponding to the (right-handed) electron isosinglet $\xi_e (= \chi_{\bar{e}})$ and neutrino isosinglet $\xi_\nu (= \chi_{\bar{\nu}})$ in the standard version. On these grounds, if t_{3L} generally denotes the weak isospin third-component eigenvalue for the case under consideration, we are consistently allowed to reproduce the leptonic electroweak Gell-Mann–Nishijima formula (with the antilepton isosinglets just in place of the right-handed lepton isosinglets) as follows:

$$-|q| = (t_{3L} + y/2)|e| \tag{5.5}$$

where $|q| = |e|$ for χ_e as well as $\chi_{\bar{e}}$, and $|q| = 0$ for χ_ν as well as $\chi_{\bar{\nu}}$. This also implies that the two fields $\chi_{\bar{e}}$ and $\chi_{\bar{\nu}}$, as far as their belonging to the $SU(2)_L$ *singlet* representation is concerned, are to be assigned the weak hypercharge eigenvalues $y = -2$ and $y = 0$, respectively. We may then *naturally* recast (5.4) in the form

$$\mathscr{L} = \mathscr{L}^{(L)} \equiv \tfrac{1}{4}[i(\overline{D}_L\gamma^\mu\partial_\mu D_L + \overline{\chi}_{\bar{e}}\gamma^\mu\partial_\mu\chi_{\bar{e}} + \overline{\chi}_{\bar{\nu}}\gamma^\mu\partial_\mu\chi_{\bar{\nu}}) + \text{H.c.}] \tag{5.6}$$

$(\overline{D}_L = D_L^\dagger\gamma^0)$ possessing manifest global invariance under the group

$$SU(2)_L \otimes U(1)_Y \tag{5.7}$$

where the subscript Y denotes a weak hypercharge variety *just pertaining* to the $SU(2)_L$ representations. The form (5.6) is still invariant under P, as applied according to (2.5), but is *no longer* invariant under P_{in}, as defined by (2.2). If local invariance is demanded as well, a gauge coupling term can be obtained:

$$\mathscr{L}_{\text{int}}^{(L)} = \tfrac{1}{2}[-g\overline{D}_L\gamma^\mu\mathbf{T}_L D_L \cdot \mathbf{W}_\mu - \tfrac{1}{2}g'(-\overline{D}_L\gamma^\mu D_L - 2\overline{\chi}_{\bar{e}}\gamma^\mu\chi_{\bar{e}})B_\mu] \tag{5.8}$$

where $\mathbf{T}_L$ is the $SU(2)_L$ generator defined in (4.4) and the $\frac{1}{2}$ factor is due to the normalized chiral field definition (1.16). The fact that the antineutrino current $\overline{\chi}_{\bar{\nu}}\gamma^\mu\chi_{\bar{\nu}}$ does *not* appear in the $U(1)_Y$ coupling sector of (5.8) is just due to the *vanishing* Y eigenvalue of $\chi_{\bar{\nu}}$. As further regards the charged gauge field $W_\mu = 2^{-1/2}(W_{1_\mu} - iW_{2_\mu})$ and its Hermitian conjugate $W_\mu^\dagger$, the former annihilates and the latter creates a (massless) vector boson with a (left-handed) weak isospin third-component eigenvalue $t_{3L} = +1$ and a conventional $|e|^2$ root $+|e|$ as assigned by (5.5) (recall that $|e|^2$ *only* can make sense at this stage). Due to the special identity $\overline{\chi}_{\bar{e}}\gamma^\mu\chi_{\bar{e}} = \overline{\xi}_e\gamma^\mu\xi_e$, we have that the *natural* coupling term (5.8) is formally coincident with what in the standard theory can but be attained *by imposing* the left-handed nature of the neutrino–electron doublet as an ad hoc external prescription. Note also the *natural* absence of antineutrino currents in the whole term (5.8) versus the *phenomenologically imposed* (equivalent) absence of right-handed neutrino currents in the ordinary coupling term.

If, rather, we choose the plus sign in (5.3)—so as to make the $|e|^2$ root *numerically* coincident with the electric charge $+e$ of the actual (massive) positron—we may then recast (5.4), quite symmetrically, as

$$\mathcal{L} = \mathcal{L}^{(R)} \equiv \tfrac{1}{4}[i(\overline{D}_R\gamma^\mu\partial_\mu D_R + \overline{\chi}_e\gamma^\mu\partial_\mu\chi_e + \overline{\chi}_\nu\gamma^\mu\partial_\mu\chi_\nu) + \text{H.c.}] \quad (5.9)$$

($\overline{D}_R = D_R^\dagger\gamma^0$). It is evident that $\mathcal{L}^{(R)}$ can be obtained from $\mathcal{L}^{(L)}$ by applying P_{in} as defined by (2.2). We are thus alternatively selecting the *right-handed*, i.e., $SU(2)_R$, variety of the weak isospin representations, where D_R is the antilepton *isodoublet* given by (4.7) and the two fields χ_e and χ_ν are now $SU(2)_R$ *singlets*. The form (5.9) possesses manifest global invariance under the group

$$SU(2)_R \otimes U(1)_{\overline{Y}} \quad (5.10)$$

where the subscript $\overline{Y}$ similarly denotes a weak hypercharge variety *just pertaining* to the $SU(2)_R$ representations, and the right-handed counterpart of formula (5.5) reads

$$+|q| = (t_{3R} + \overline{y}/2)|e| \quad (5.11)$$

thus particularly implying $\overline{y} = +2$ for χ_e and $\overline{y} = 0$ for χ_ν. The condition of local invariance under (5.10) requires $\mathcal{L}^{(R)}$ to be supplemented by a coupling term like

$$\mathcal{L}^{(R)}_{\text{int}} = \tfrac{1}{2}[-g\overline{D}_R\gamma^\mu\mathbf{T}_R D_R \cdot \mathbf{W}_\mu - \tfrac{1}{2}g'(+\overline{D}_R\gamma^\mu D_R + 2\overline{\chi}_e\gamma^\mu\chi_e)B_\mu] \quad (5.12)$$

$\mathbf{T}_R$ is the $SU(2)_R$ generator defined in (4.4). Since $(-\boldsymbol{\tau})/2$, rather than $\boldsymbol{\tau}/2$, is the *effective* SU(2) generator resulting from the application of $\mathbf{T}_R$ to D_R, the two charged gauge fields W_μ and $W_\mu^\dagger$ are now (the former) annihilating and (the latter) creating a (massless) vector boson with a (right-handed) weak-isospin third-component eigenvalue $t_{3R} = -1$ and a corresponding $|e|^2$ root $-|e|$ as assigned by (5.11).

Actually, both coupling terms (5.8) and (5.12) come from two *P-invariant* free Lagrangian densities, as can be checked by applying P to either Eq. (5.6) or Eq. (5.9) according to (2.5); and they may *themselves* be individually P-invariant, provided the gauge fields $\mathbf{W}_\mu$ and B_μ are assumed to be *axial* vectors in space-time. This is because the left- and right-handed weak isospin currents, as covariantly recast in terms of $\mathbf{T}^W$ with the help of Eqs. (4.2) and (4.4),

$$\overline{D}_L\gamma^\mu\mathbf{T}_L D_L = \overline{D}_L\gamma^\mu\mathbf{T}^W D_L, \qquad \overline{D}_R\gamma^\mu\mathbf{T}_R D_R = \overline{D}_R\gamma^\mu\mathbf{T}^W D_R \quad (5.13)$$

already behave as axial vectors under

$$D_{L,R} \to \gamma^0 D_{L,R}, \quad \overline{D}_{L,R} \to \overline{D}_{L,R}\gamma^0 \quad (5.14)$$

and the same clearly holds for the related weak hypercharge currents (if y

and $\bar{y}$ are consistently taken as pseudoscalars and are turned into $-y$ and $-\bar{y}$). On the other hand, it is also to be said that Eqs. (5.5) and (5.11) cannot strictly have invariant forms under P, since either root $\mp|q|$ stands by itself for a scalar (and not a pseudoscalar) quantity: the $\mp|e|$ current (5.3) behaves as a (true) vector in space-time and is left unvaried by pseudoscalar-charge conjugation P_{in}, as acting according to (2.2) or (2.9). That, however (in the light of the conclusions drawn in Section 2), implies no actual failure of P symmetry, by virtue of the fact that $|q|^2$ *only* can really make sense for a chiral field (unless, of course, $|q|^2 = 0$). On applying P (or P_{in}) we may well understand, therefore, that we are also *conventionally inverting* the $|q|^2$ root so as to overcome the formal impasse above. This can be expressed by saying that P (or P_{in}) *alone is here acting the same as CP* (or CP_{in}), where, by definition,

$$C: \quad -|q| \rightleftarrows +|q| \tag{5.15}$$

and where invariance under C just means that we could have equally started by setting, in place of Eqs. (5.5) and (5.11), a pair of relations with interchanged $|q|^2$ roots,

$$+|q| = -(t_{3L} + y/2)|e|, \qquad -|q| = -(t_{3R} + \bar{y}/2)|e| \tag{5.16}$$

In spite of their established P-invariance property, the single coupling terms (5.8) and (5.12) are not invariant with respect to P_{in}, just as happens for the P_{in}-*conjugated* free terms (5.6) and (5.9). If, for instance, we apply both (2.2) and (4.1) to Eq. (5.8), we obtain

$$P_{\mathrm{in}}: \quad \begin{array}{l} \overline{D}_L\gamma^\mu \mathbf{T}^W D_L \to \overline{D}_R\gamma^\mu(-\mathbf{T}^W)D_R \\ (-\overline{D}_L\gamma^\mu D_L - 2\overline{\chi}_{\bar{e}}\gamma^\mu x_{\bar{e}}) \to -(\overline{D}_R\gamma^\mu D_R + 2\,\overline{\chi}_{e}\gamma^\mu\chi_{e}) \end{array} \tag{5.17}$$

and we see that *Eq. (5.12) may similarly be the* P_{in}*-conjugated* (or CP_{in}-conjugated) *counterpart of Eq. (5.8), provided*

$$P_{\mathrm{in}}(CP_{\mathrm{in}}): \quad \mathbf{W}_\mu \to -\mathbf{W}_\mu, \quad B_\mu \to -B_\mu \tag{5.18}$$

Such a condition, on the other hand, is exactly what should be expected for *axial* vectors in space-time (see the end of Section 1) and turns out to legitimate the formal use of the *same* gauge fields (and coupling constants) in (5.8) as in (5.12). From applying both (2.2) and (4.1) as above, it also follows that P_{in} does not cause any net change in the effective isospin $\frac{1}{2}\boldsymbol{\tau}$ and may then be interpreted, in view of (5.18), as representing the whole inversion operation in the corresponding effective isospin space, with the P_{in} (CP_{in}) parity of $\mathbf{W}_\mu$ amounting to a “G parity.”[28] Yet it should be noticed that the isodoublet $D_L^{(P_{\mathrm{in}})}$ (obtained from D_L by *just* making the substitutions $\chi_\nu \to \chi_{\bar{\nu}}$, $\chi_e \to \chi_{\bar{e}}$) does not look like the “G conjugate” of D_L. This is explained by the fact that $D_L^{(P_{\text{in}})}$ is acted upon by the *new* weak isospin operator $\mathbf{T}^{W(P_{\text{in}})} = \frac{1}{2}\boldsymbol{\tau}\gamma^5 = -\mathbf{T}^{W}$ (and not by $\mathbf{T}^{W}$): so, though it contains antilepton fields, it needs no internal changes to belong to the same effective SU(2) representation as D_L (more precisely, the required sign inversion of the weak isospin third-component eigenvalues is already ensured by the transformation $\mathbf{T}^{W} \to \mathbf{T}^{W(P_{\text{in}})}$).

Before going further, it is worth pointing out the following peculiar feature of the present formulation, strictly connected with the P_{in} property of interchanging left- and right-handed chiral fields: in building either of the Lagrangian densities (5.6) and (5.9) from the original free Lagrangian density (5.4), we have actually brought about a *spontaneous breaking* of P_{in} symmetry, which has led us to an either "left-handed" or "right-handed" approach, based on gauge invariance with respect to either the group (5.7) or the group (5.10). A similar remark may also apply to C invariance above, which, in either case, may be said to be *spontaneously broken* by the corresponding choice of the single $|e|^2$ root to be assigned to both (electron and positron) chiral fields χ_{e} and$\chi_{\bar{\text{e}}}$.

6. A UNIFIED ($L + R$)-COVARIANT FORMALISM

In the P-symmetric framework under consideration, the two covariant gauge couplings (5.8) and (5.12) are obviously equivalent and *alternative*. Their derivation merely depends on whether the negative or positive $|e|^2$ root is being chosen to mark *both* chiral fields χ_{e} and $\chi_{\bar{\text{e}}}$: according to the choice made, the $\text{SU}(2)_L \otimes \text{U}(1)_Y$ or $\text{SU}(2)_R \otimes \text{U}(1)_{\bar{Y}}$ symmetry group is in turn selected.

Unlike χ_ν and $\chi_{\bar{\nu}}$, the two fields χ_{e} and $\chi_{\bar{\text{e}}}$ are *together* present in both (5.8) and (5.12). So, if we want to build an overall $L + R$ formalism that may *unambiguously* merge such couplings, we should also find some *explicit* way of marking the *orthogonality* between the isomultiplets pertaining to $\text{SU}(2)_L \otimes \text{U}(1)_Y$ representations and those pertaining to $\text{SU}(2)_R \otimes \text{U}(1)_{\bar{Y}}$ representations.

For the purpose in question, the two roots $\mp|q\,|$ ($|q|_{\nu,\bar{\nu}} = 0$, $|q|_{\text{e},\bar{\text{e}}} = |e|$) may be associated to a pair of *Casimir operators*,

$$\mathcal{P}_- \equiv |-\rangle\langle-|, \qquad \mathcal{P}_+ \equiv |+\rangle\langle+| \tag{6.1}$$

$$\mathcal{P}_- + \mathcal{P}_+ = 1, \qquad \mathcal{P}_-\mathcal{P}_+ = \mathcal{P}_+\mathcal{P}_- = 0, \qquad \mathcal{P}_-^2 = \mathcal{P}_-, \qquad \mathcal{P}_+^2 = \mathcal{P}_+$$

which should in turn *label* the two sets of $\text{SU}(2)_L \otimes \text{U}(1)_Y$ and $\text{SU}(2)_R \otimes \text{U}(1)_{\bar{Y}}$ representations and should just enable one to *select* either of them. We can thus start anew with the free Lagrangian density (5.4) and recast it in the "dressed" form

$$\mathscr{L} = \mathscr{L}^{(L+R)} \equiv \mathscr{L}^{(L)}\mathscr{P}_- + \mathscr{L}^{(R)}\mathscr{P}_+ \tag{6.2}$$

$\mathscr{L}^{(L)}$ and $\mathscr{L}^{(R)}$ being given by (5.6) and (5.9). Doing this, we have reintroduced the *spontaneous breaking* of P_{in} symmetry, but we have also *explicitly* connected both $\mathscr{L}^{(L)}$ with the root $-|q|$ and $\mathscr{L}^{(R)}$ with the root $+|q|$. That implies the definition of an overall "dressed" covariant Fock space being the sum of *two* Fock-space varieties: the former variety, fit for $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$, is identified by the "label" $|-\rangle$ (in ket notation) or $\langle -|$ (in bra notation) and the latter variety, fit for $\mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{\bar{Y}}$, is similarly identified by $|+\rangle$ or $\langle +|$. The relations (5.5) and (5.11), the former coupled to $\mathscr{L}^{(L)}$ and the latter to $\mathscr{L}^{(R)}$, should accordingly be replaced by the "dressed" ones

$$-|q|\mathscr{P}_- = (t_{3_L} + y/2)|e|\mathscr{P}_- \tag{6.3}$$

and

$$+|q|\mathscr{P}_+ = (t_{3_R} + \bar{y}/2)\,|e|\mathscr{P}_+ \tag{6.4}$$

Such formulas can be merged into the single formula

$$-\,|q|(\mathscr{P}_- - \mathscr{P}_+) = [(t_{3_L} + y/2)\mathscr{P}_- + (t_{3_R} + \bar{y}/2)\mathscr{P}_+]|e| \tag{6.5}$$

which refers to (6.2) as a whole; this should be related to the fact that the Lagrangian density (6.2) may be said to possess manifest global invariance under an overall gauge group defined as

$$\begin{aligned} &[\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y] \oplus [\mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{\bar{Y}}] \\ &\quad \equiv [\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y]\,\mathscr{P}_- + [\mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{\bar{Y}}]\mathscr{P}_+ \end{aligned} \tag{6.6}$$

If local invariance with respect to the group (6.6) is demanded, a *double*-structured coupling term like

$$\mathscr{L}_{\text{int}}^{(L+R)} = \mathscr{L}_{\text{int}}^{(L)}\mathscr{P}_- + \mathscr{L}_{\text{int}}^{(R)}\mathscr{P}_+ \tag{6.7}$$

is found, where $\mathscr{L}_{\text{int}}^{(L)}$ and $\mathscr{L}_{\text{int}}^{(R)}$ are given by (5.8) and (5.12).

A glance at (6.5) shows that $-|q| \rightleftarrows +|q|$ can now be equivalently obtained as a result of $\mathscr{P}_- \rightleftarrows \mathscr{P}_+$. One may then put

$$C|-\rangle = |+\rangle, \qquad C|+\rangle = |-\rangle \tag{6.8}$$

and write, in place of (5.15),

$$C\mathscr{P}_-C^{-1} = \mathscr{P}_+, \qquad C\mathscr{P}_+C^{-1} = \mathscr{P}_- \tag{6.9}$$

($C^{-1} = C^\dagger = C$). Neither of the forms (6.2), (6.7) is clearly invariant under (6.9): on passing to the $L + R$ formalism in hand, the *spontaneous breaking* of C symmetry becomes manifest. Such an outcome corresponds to the fact that the *vacuum* Fock state is no longer C-invariant, but may in turn take either of the C-conjugated dressed forms $|0\rangle|-\rangle$ or $|0\rangle|+\rangle$, according to whether one is referring to an $SU(2)_L \otimes U(1)_Y$ or $SU(2)_R \otimes U(1)_{\bar{Y}}$ representation, respectively (this, of course, does not imply an electrically charged vacuum, since $|0\rangle$ is still a no-particle state in covariant terms). The overall result is that both forms (6.2) and (6.7) can at most exhibit CP_{in} invariance (besides P invariance), in spite of the fact that P_{in} *alone* is already able to provide for $\mathcal{L}^{(L)} \rightleftarrows \mathcal{L}^{(R)}$ as well as $\mathcal{L}^{(L)}_{\text{int}} \rightleftarrows \mathcal{L}^{(R)}_{\text{int}}$ and is then really a symmetry operation (like C): as generally happens for a symmetry that is spontaneously broken, both P_{in} and C individual symmetries of the original Lagrangian density (5.4) have only been "hidden"[29] on recasting it in the form (6.2).

7. AN INNER MECHANISM OF MASS GENERATION AS AN ESSENTIAL REQUIREMENT FOR A CHARGED LEPTON TO BECOME AN ACTUAL EIGENSTATE OF ELECTRIC CHARGE

Within the above $L + R$ formal description we may build from (5.3) a "dressed"conserved electric current operator that reads as follows:

$$\mathcal{J}^{(Q)} \equiv \tfrac{1}{2}[-|q|(\bar{\chi}_e\gamma^\mu\chi_e + \bar{\chi}_{\bar{e}}\gamma^\mu\chi_{\bar{e}})\mathcal{P}_- + |q|(\bar{\chi}_e\gamma^\mu\chi_e + \bar{\chi}_{\bar{e}}\gamma^\mu\chi_{\bar{e}})\mathcal{P}_+] \qquad (7.1)$$

where $|q| = |e|$. If we make the substitution

$$\chi_e = 2^{-1/2}(\psi_e - \psi_{\bar{e}}), \qquad \chi_{\bar{e}} = 2^{-1/2}(\psi_e + \psi_{\bar{e}}) \qquad (\psi_{\bar{e}} = \gamma^5\psi_e) \quad (7.2)$$

(together with its adjoint) in (7.1), we see that this current becomes strongly reminiscent of the massive scalar-charge free current (1.37), with $\mathcal{P}_-$ and $\mathcal{P}_+$ just corresponding to $\mathcal{P}_f$ and $\mathcal{P}_{\bar{f}}$. A glance at (7.1) directly shows, however, what has been pointed out in Section 3: as long as χ_e and $\chi_{\bar{e}}$ are *massless*, so that the two Dirac combinations

$$\psi_e = 2^{-1/2}(\chi_e + \chi_{\bar{e}}), \qquad \psi_{\bar{e}} = 2^{-1/2}(-\chi_e + \chi_{\bar{e}}) \qquad (7.3)$$

may only be *mixtures* of them, the electron and positron *are not allowed* to look like actual (conjugated) eigenstates of electric charge. On the other hand, as soon as mass is acquired, both ψ_e and $\psi_{\bar{e}}$ *stop* being bound to be mixtures of χ_e and $\chi_{\bar{e}}$, and may really become two (conjugated) electric charge eigenfields, each with one independent nonmeasurable phase. Here the existence of an electric charge *superselection rule* thus becomes a *non*trivial requirement, which can indeed be fulfilled only if the electron and positron are made *massive*.

Hence, strictly speaking, the appearance of electron and positron masses should now be expected to have just an *internal* origin, connected with the *emerging* dynamical nature of electric charge, and no longer a purely "accidental," *external* origin, due to an ad hoc prescribed coupling to the Higgs boson. But how is it possible to think of any SSB mechanism that is able to make the electron and positron massive *apart from* the Higgs boson existence? What would be required is some "effective" Higgs-like isodoublet field *not* being a true external field: *its action upon the electron (positron) should just be the (internal) one of giving mass, without also involving the coupling to an outside real particle* (like the Higgs boson).

Take for a moment the Lagrangian density relevant to the usual Higgs isodoublet field:

$$\mathscr{L}_{\mathrm{Higgs}} = \partial^\mu \phi^\dagger \partial_\mu \phi - a\phi^\dagger\phi - b(\phi^\dagger\phi)^2 \qquad (a < 0,\, b > 0) \tag{7.4}$$

where

$$2\langle \phi^\dagger \phi \rangle_0 = |a|/b \equiv v^2 \tag{7.5}$$

If we make the two coefficients a, b *identically go to zero in modulus*, we can clearly keep the vacuum expectation value (7.5) *unvaried*. Hence, as an extrapolation of (7.4), one may *also* think of an isodoublet field $\phi(x)$ being *massless* and further retaining a *non*zero vacuum expectation value under the strict condition of *absence* of any quartic self-coupling term:

$$\mathscr{L}_{\mathrm{Higgs}} = \partial^\mu \phi^\dagger \partial_\mu \phi, \qquad 2\langle \phi^\dagger \phi \rangle_0 = v^2 \neq 0 \tag{7.6}$$

Let us consider (7.6) in place of (7.4). In the model so introduced, the existence of the Higgs boson massive excitation is naturally prevented, and one may further redefine $\phi(x)$ anew as a field being endowed with *three* (massless) freedom degrees only. According to whether we are dealing with the $SU(2)_L \otimes U(1)_Y$ or $SU(2)_R \otimes U(1)_{\overline{Y}}$ representation—$\frac{1}{2}\boldsymbol{\tau}$ and $\frac{1}{2}(-\boldsymbol{\tau})$ being the effective SU(2) generators in the two cases—we may express $\phi(x)$ in either of the equivalent general isodoublet forms

$$\begin{aligned} \phi(x) &\equiv \phi_L(x) = \exp[i\boldsymbol{\theta}(x) \cdot \tfrac{1}{2}\boldsymbol{\tau}]\phi_0, \\ \phi(x) &\equiv \phi_R(x) = \exp[-i\boldsymbol{\theta}(x) \cdot \tfrac{1}{2}(-\boldsymbol{\tau})]\phi_0 \end{aligned} \tag{7.7}$$

where

$$\phi_0 = \langle \phi(x) \rangle_0 = 2^{-1/2}\begin{pmatrix} 0 \\ v \end{pmatrix} \tag{7.8}$$

and where $\boldsymbol{\theta}(x)$ is a Goldstone axial-isovector field. It is also to be noted that according to whether $\phi(x)$ is being classified within the $SU(2)_L \otimes U(1)_Y$ or $SU(2)_R \otimes U(1)_{\overline{Y}}$ representation, the vacuum expectation value (7.8) will be in turn associated with an eigenvalue $y = +1$ of the hypercharge Y or an eigenvalue $\overline{y} = -1$ of the hypercharge $\overline{Y}$. Under the local gauge invariance condition, the three components of $\boldsymbol{\theta}(x)$ can clearly be eliminated by a standard procedure to give mass to three of the four gauge fields involved: for instance, the $[SU(2)_L \otimes U(1)_Y]$ locally invariant Lagrangian density generalizing (7.6) reads (omitting the terms for the gauge fields alone)

$$\{[\partial^\mu + i(g/2)\boldsymbol{\tau}\cdot\mathbf{W}^\mu + i(g'/2)B^\mu]\phi_L\}^\dagger[\partial_\mu + i(g/2)\boldsymbol{\tau}\cdot\mathbf{W}_\mu + i(g'/2)B_\mu]\phi_L$$

$$= \{[\partial^\mu + i(g/2)\boldsymbol{\tau}\cdot\mathbf{W}'^\mu + i(g'/2)B^\mu]\phi_0\}^\dagger[\partial_\mu + i(g/2)\boldsymbol{\tau}\cdot\mathbf{W}'_\mu + i(g'/2)B_\mu]\phi_0 \quad (7.9)$$

with $\mathbf{W}'^\mu = \mathbf{W}^\mu + g^{-1}\partial^\mu\boldsymbol{\theta}(x)$, and the $SU(2)_R \otimes U(1)_{\overline{Y}}$ counterpart of (7.9) can be obtained by the replacements $i \to -i$, $\boldsymbol{\tau} \to -\boldsymbol{\tau}$, and $\phi_L \to \phi_R$. The main difference from the standard Higgs isodoublet model applied to the electroweak theory[24,25] lies therefore in the *absence* of the Higgs boson excitation (and related coupling terms). For this reason, the field $\phi(x)$ under consideration may simply be referred to as a *"vacuum" Higgs field*. Its *effective* presence (being such as to preserve the gauge symmetry) should *just* express the mass-generating power that the electric charge is here required to exert in order to be able to yield the electromagnetic gauge coupling.

In strict analogy to the usual procedure, the leptonic Lagrangian mass term (whose appearance is now essential) should be obtained as a result of a symmetry-preserving coupling to *what is left* of $\phi(x)$ after the elimination of $\boldsymbol{\theta}(x)$. In the overall $L + R$ picture, it takes the form of a *double*-structured term

$$\mathcal{L}_{\text{mass}} = \tfrac{1}{2}G_e[(\overline{D}_L\chi_{\bar{e}}\mathcal{P}_- + \overline{D}_R\chi_e\mathcal{P}_+)\phi_0 + \text{H.c.}] \quad (7.10)$$

By substituting (7.8) it follows that the constant G_e multiplied by $2^{-1/2}v$ gives the actual modulus $|m_e|$ of either the electron or positron proper mass:

$$\begin{aligned}\mathcal{L}_{\text{mass}} &= 2^{-1/2}G_e v[\tfrac{1}{2}(\overline{\chi}_e\chi_{\bar{e}} + \overline{\chi}_{\bar{e}}\chi_e)\mathcal{P}_- + \tfrac{1}{2}(\overline{\chi}_e\chi_{\bar{e}} + \overline{\chi}_{\bar{e}}\chi_e)\mathcal{P}_+] \\ &= |m_e|(\overline{\psi}_e\psi_e\mathcal{P}_- - \overline{\psi}_{\bar{e}}\psi_{\bar{e}}\mathcal{P}_+) \quad (7.11)\end{aligned}$$

As can easily be checked, this is a mass term just amounting to that included in (1.28) (note, in particular, the opposite sign of the electron and positron proper mass values). Here G_e is no longer fixed by the interaction with an outside real particle (such as the Higgs boson): it should merely be taken as a "vacuum coupling constant" (pertaining to the special charged lepton variety under consideration). Of course, condition (7.8) ensures ϕ_0 to have a null electric charge value, so as to be left invariant by the U(1) gauge group that the electric charge is now enabled to generate as a subgroup of (6.6).

As is implicitly shown by (7.9), the acquired masses of the charged gauge fields W_μ, $W_\mu^\dagger$ and of the final neutral gauge field $Z_\mu = \cos\theta_W W_{3\mu} - \sin\theta_W B_\mu$ (θ_W is the Weinberg angle) are exactly the same as in the standard formulation. Note, on the other hand, that such a SSB mechanism is left–right (or P_{in}) symmetric and cannot affect the already established *axial* vector nature of either W_μ, $W_\mu^\dagger$ or $W_{3\mu}$, B_μ. Thus we may finally replace (5.18) with

$$P_{\text{in}}(CP_{\text{in}}): \quad W_\mu, W_\mu^\dagger, Z_\mu \to -W_\mu, -W_\mu^\dagger, -Z_\mu \tag{7.12}$$

(from now on, W_μ and $W_\mu^\dagger$ are directly understood to be massive, without the use of a primed notation). The fact that W_μ and $W_\mu^\dagger$ themselves are C-invariant gauge fields (C being the considered covariant operation of scalar charge reversal) should not be surprising. In the present formalism the well-known "*CP* symmetry" of chiral processes just amounts to a pure P symmetry, as if all scalar charges involved in those processes are momentarily subject to a *maximal sign uncertainty*: in particular, as can be drawn from the concluding remarks of Section 5, the effective transformation interchanging W_μ and $W_\mu^\dagger$ is merely given by $P_{\text{in}} \exp[-i(\pm\tau_2/2)\pi]$ [the plus or minus sign applying respectively to the $SU(2)_L$ case, with an effective isospin $\boldsymbol{\tau}/2$, or the $SU(2)_R$ one, with an effective isospin $-\boldsymbol{\tau}/2$]. On the other hand, it should also be recalled that C, in line with its covariant nature, does *not* affect annihilation and creation operators. If the standard Fourier expansions of W_μ and $W_\mu^\dagger$ are taken into account—the former including "antiparticle"-creation (as well as "particle"-annihilation) operators and the latter including "antiparticle"-annihilation (as well as "particle"-creation) operators—then a *non*covariant operation of charge conjugation (available for *positive* energies only) can also be introduced: it just interchanges "particle" and "antiparticle" and just amounts to $P_{\text{in}} \exp[-i(\pm\tau_2/2)\pi]$.

8. NEW PECULIAR FEATURES OF THE FINAL LEPTONIC ELECTROWEAK COUPLING

As we have seen above, what is left of the effective "vacuum" Higgs field $\phi(x)$ after the "absorption" of the three Goldstone degrees of freedom is just its (nonzero) vacuum expectation value ϕ_0. The coupling of ϕ_0 to the lepton–antilepton system makes two mere (original) chiral-field mixtures like those in (7.3) become the *actual* (Dirac) proper-mass eigenfields (with opposite eigenvalues) of the physical electron and positron. This statement can be reversed by saying that, as a result of the coupling to ϕ_0, the electron and positron proper-mass eigenfields are to be identified now with $\psi_e(x)$ and $\psi_{\bar{e}}(x)$, and *no longer* with $\chi_e(x)$ and $\chi_{\bar{e}}(x)$. With regard to this, the individual presence of the chiral fields $\chi_e(x)$ and $\chi_{\bar{e}}(x)$ in the weak isospin gauge couplings may really survive only as a *dynamical* constraint due to the pseudoscalar nature of the weak isospin (in its single components): such a nature, expressed by (4.2), still involves a *double* variety of SU(2) representations, marked by the two *superselecting* fundamental generators $\mathbf{T}_L = \mathbf{T}^W X_L$ and $\mathbf{T}_R = \mathbf{T}^W X_R$, with $X_L = \frac{1}{2}(1 - \gamma^5)$ and $X_R = \frac{1}{2}(1 + \gamma^5)$.

The appearance of the electron–positron mass term (7.11) makes it sound to introduce a one-particle electric charge "operator" formally as, in view of Eq. (6.5),

$$Q(|q|) = -|q|\,(\mathcal{P}_- - \mathcal{P}_+) \qquad (8.1)$$
$$\equiv -|q|L_e \qquad (|q|_{\nu,\bar{\nu}} = 0,\ |q|_{e,\bar{e}} = |e|)$$

where $L_e = (+1)\mathcal{P}_- + (-1)\mathcal{P}_+$ may accordingly be taken as being a one-particle electron–lepton number "operator" with (particle and antiparticle) eigenvalues $l_e = (+1)$ and $\bar{l}_e = (-1)$. Charge (8.1) is clearly defined as a scalar quantity (relative to ordinary space inversion). It then commutes with P_{in}, just like the general scalar-charge operator (1.35); furthermore, it anticommutes with C as given by (6.9), so that C itself may at last be strictly interpreted as a (covariant) scalar-charge conjugation operator. Another remarkable property of charge (8.1) is that it applies to the *whole* space as the sum of the two (C-conjugated) Fock-space varieties marked (in ket notation) by $|-\rangle$ and $|+\rangle$; this actually shows that the U(1) gauge group generated by $Q(|q|)$ can only be a subgroup of the *overall* group (6.6). The electric charge so defined will clearly generate a neutral gauge field A_μ behaving as a (true) *vector* and then commuting with P_{in} (according to the covariant rule established at the end of Section 1). Such an outcome is consistent with the fact that A_μ is not originated by (7.9) and is only required to be orthogonal to Z_μ. As a consequence, a link like $A_\mu = \cos\theta_W\, W_{3\mu} + \sin\theta_W\, B_\mu$ should strictly be taken as being *neither* P_{in}- *nor* P-invariant (since both $W_{3\mu}$ and B_μ anticommute with P_{in} and are axial vectors).

In view of this, the two alternative weak hypercharge varieties $Y(y) = y\mathcal{P}_-$ and $\bar{Y}(\bar{y}) = \bar{y}\mathcal{P}_+$ are to be recast in terms of $Q = Q(|q|)$ and to be "split" as follows:

$$\begin{aligned} Y(t_{3L}, |q|) &= 2(-t_{3L} - |q|/|e|)\mathcal{P}_-, \\ \bar{Y}(t_{3R}, |q|) &= 2(-t_{3R} + |q|/|e|)\mathcal{P}_+ \end{aligned} \qquad (8.2)$$

Thus, if we substitute (8.2) into (6.5) and make an explicit allowance for W_μ, $W^\dagger_\mu$, Z_μ, and A_μ, we come to the final overall $L + R$ coupling, which formally joins two *equivalent* (and alternative) covariant electroweak couplings in terms of either *leptons* or *antileptons* (with both positive and negative energies). It reads

$$\mathcal{L}_{int}(e, \nu; \bar{e}, \bar{\nu}) = [\mathcal{L}_{int}(e, \nu) + |e|\bar{\psi}_e\gamma^\mu\psi_e A_\mu]\mathcal{P}_-$$
$$+ [\mathcal{L}_{int}(\bar{e}, \bar{\nu}) - |e|\bar{\psi}_{\bar{e}}\gamma^\mu\psi_{\bar{e}} A_\mu]\mathcal{P}_+ \qquad (8.3)$$

($g \sin\theta_W = g' \cos\theta_W = |e|$), where (directly passing to a manifestly covariant form)

$$\mathcal{L}_{\rm int}(\mathrm{e},\,\nu) = -2^{-3/2}g(\overline{D}_L\gamma^\mu T^{\mathrm{w}+}D_L W_\mu + \overline{D}_L\gamma^\mu T^{\mathrm{w}-}D_L W_\mu^\dagger) \\ - (g/2\cos\theta_{\mathrm{W}})(\overline{D}_L\gamma^\mu T_3^{\mathrm{W}}D_L + 2\sin^2\theta_{\mathrm{W}}\overline{\psi}_{\mathrm{e}}\gamma^\mu\psi_{\mathrm{e}})Z_\mu \tag{8.4}$$

$$\mathcal{L}_{\rm int}(\bar{\mathrm{e}},\,\bar{\nu}) = -2^{-3/2}g(\overline{D}_R\gamma^\mu T^{\mathrm{W}+}D_R W_\mu + \overline{D}_R\,\gamma^\mu T^{\mathrm{W}-}D_R W_\mu^\dagger) \\ - (g/2\cos\theta_{\mathrm{W}})(\overline{D}_R\gamma^\mu T_3^{\mathrm{W}}D_R - 2\sin^2\theta_{\mathrm{W}}\overline{\psi}_{\bar{\mathrm{e}}}\gamma^\mu\psi_{\bar{\mathrm{e}}})Z_\mu \tag{8.5}$$

$[T^{\mathrm{W}\pm} = (T_1^{\mathrm{W}} \pm iT_2^{\mathrm{W}})]$, and where, in line with Eqs. (1.31) and (7.2) and by use of Eqs. (6.1), one can write

$$\overline{\psi}_{\mathrm{e}}\gamma^\mu\psi_{\mathrm{e}} = \overline{\psi}_{\bar{\mathrm{e}}}\gamma^\mu\psi_{\bar{\mathrm{e}}} = \tfrac{1}{2}(\overline{\chi}_{\mathrm{e}}\gamma^\mu\chi_{\mathrm{e}} + \overline{\chi}_{\bar{\mathrm{e}}}\gamma^\mu\chi_{\bar{\mathrm{e}}}) \\ = \tfrac{1}{2}[\overline{\Psi}(\mathrm{e},\bar{\mathrm{e}})\gamma^\mu\Psi(\mathrm{e},\bar{\mathrm{e}}) + \overline{\Psi}^{(P_{\rm in})}(\mathrm{e},\bar{\mathrm{e}})\gamma^\mu\Psi^{(P_{\rm in})}(\mathrm{e},\bar{\mathrm{e}})] \\ \equiv J(\mathrm{e},\bar{\mathrm{e}}) \tag{8.6}$$

$$\begin{aligned} \Psi(\mathrm{e},\bar{\mathrm{e}}) &\equiv \psi_{\mathrm{e}}\langle -| + \psi_{\bar{\mathrm{e}}}\langle +|, \qquad \overline{\Psi}(\mathrm{e},\bar{\mathrm{e}}) \equiv \overline{\psi}_{\mathrm{e}}|-\rangle + \overline{\psi}_{\bar{\mathrm{e}}}|+\rangle \\ \Psi^{(P_{\rm in})}(\mathrm{e},\bar{\mathrm{e}}) &\equiv \psi_{\mathrm{e}}\langle -| - \psi_{\bar{\mathrm{e}}}\langle +|, \qquad \overline{\Psi}^{(P_{\rm in})}(\mathrm{e},\bar{\mathrm{e}}) \equiv \overline{\psi}_{\mathrm{e}}|-\rangle - \overline{\psi}_{\bar{\mathrm{e}}}|+\rangle \end{aligned} \tag{8.7}$$

In strict analogy with the massless case, the sign of both electric charge and weak-isospin third-component of the boson being annihilated by W_μ or created by $W_\mu^\dagger$ is not *a priori* fixed, but consistently depends on whether W_μ and $W_\mu^\dagger$ are taken within the $\mathrm{SU}(2)_L$ scheme, as in (8.4), or the $\mathrm{SU}(2)_R$ one, as in (8.5).

The formalism introduced by (8.6) and (8.7) is just an exemplification of the generalized field formalism (for massive spin-1/2 point fermions and antifermions) mentioned in Section 1. The "bare" current term $J(\mathrm{e}, \bar{\mathrm{e}})$ acts in a single, strictly covariant, Fock space, $\mathcal{F}^\circ(\mathrm{e}, \bar{\mathrm{e}})$, equally available for either (positive- and negative-energy) electrons or (positive- and negative-energy) positrons. By taking the direct product of $\mathcal{F}^\circ(\mathrm{e}, \bar{\mathrm{e}})$ with the (two-dimensional) internal space $\mathcal{S}_{\rm in}(l_{\mathrm{e}}, \bar{l}_{\mathrm{e}})$ spanned by the "dressing" vector basis $(|-\rangle, |+\rangle)$, one can build from $\mathcal{F}^\circ(\mathrm{e}, \bar{\mathrm{e}})$ a "dressed" covariant Fock space

$$\mathcal{F}(\mathrm{e},\bar{\mathrm{e}}) \equiv \mathcal{F}^\circ(\mathrm{e},\bar{\mathrm{e}}) \otimes \mathcal{S}_{\rm in}(l_{\mathrm{e}},\bar{l}_{\mathrm{e}}) \tag{8.8}$$

This allows an actual *doubling* of $\mathcal{F}^\circ(\mathrm{e}, \bar{\mathrm{e}})$ into a strict *electronic* and a strict *positronic* covariant Fock space, the latter being the C-conjugated counterpart of the former, and C always being the *covariant* scalar-charge conjugation given by (6.9). Accordingly, a "dressed" (doubled) electromagnetic current can be formally built from $J(\mathrm{e}, \bar{\mathrm{e}})$,

$$\mathcal{J}^{(Q)}(\mathrm{e},\bar{\mathrm{e}}) = QJ(\mathrm{e},\bar{\mathrm{e}}) = J(\mathrm{e},\bar{\mathrm{e}})Q = -|e|[J(\mathrm{e},\bar{\mathrm{e}})\mathcal{P}_- - J(\mathrm{e},\bar{\mathrm{e}})\mathcal{P}_+] \tag{8.9}$$

which is included in (8.3) as the overall one coupled to A_μ. This current, being such that

$$C\mathcal{J}^{(Q)}(\mathrm{e}, \bar{\mathrm{e}}) = -\mathcal{J}^{(Q)}(\mathrm{e}, \bar{\mathrm{e}})C \tag{8.10}$$

is fit for acting in the Fock space (8.8), and the same remarks apply to it as to the general scalar-charge current (1.37).

In the *leptonic* covariant picture given by (8.4), the actual electron has a *twofold* (Dirac and chiral) field representation, $\psi_e(x)$ and $\chi_e(x)$, and the same can be said for the actual positron in the (alternative) *antileptonic* picture given by (8.5): this is intimately related to the different (scalar and pseudoscalar) natures of electric charge, on one hand, and weak-isospin charges, on the other. The fact that the whole leptonic (antileptonic) term in (8.3) is coupled to $\mathcal{P}_-$ ($\mathcal{P}_+$) directly follows from the manifest *spontaneous breaking* of C symmetry (see Section 6). In both (leptonic and antileptonic) pictures, no physical trace is left of the original, either *positron* or *electron*, isosinglet current, which turns out to be formally embodied in the final electromagnetic current term, either $-|e|\bar{\psi}_e\gamma^\mu\psi_e\mathcal{P}_-$ or $+|e|\bar{\psi}_{\bar{e}}\gamma^\mu\psi_{\bar{e}}\mathcal{P}_+$, according to (8.6): such a disappearance is essential to understanding how, in each single picture, what was originally a given $|e|^2$ root, at most expressing (no matter for the sign) the mere electric charge *magnitude* common to both involved (massless) chiral electron and positron, may now fully express the actual electric charge value of the *only* (massive) Dirac electron or positron. It should be emphasized, however, that as in the massless case, each of the two equivalent net weak coupling terms (8.4) and (8.5) can be obtained from the other by just applying P_{in} according to (2.2), (4.1), and (7.12): so, an actual lepton (antilepton) involved in a weak coupling still behaves apparently like a *pseudoscalar*-charge eigenstate that may at most be carrying also some scalar charges *definite only in magnitude*. This holds even for electrically charged leptons (antileptons), which are then subject to a charge *dualism*: in the electromagnetic coupling, of course, they just on the contrary look like *scalar*-charge eigenstates with pseudoscalar charges quite indefinite in sign.

As particularly regards the charged weak current couplings included in either (8.4) or (8.5), they can individually be seen to have a true *P-invariant* form as in the (massless) starting version. The crucial point lies in the fact that herein the Dirac parity operator γ^0 applies *directly* to chiral fields, e.g., one has $(1 - \gamma^5)\psi \rightarrow \gamma^0(1 - \gamma^5)\psi = (1 + \gamma^5)\gamma^0\psi$ *instead of* $(1 - \gamma^5)\psi \rightarrow (1 - \gamma^5)\gamma^0\psi$, so that a P-conjugated matrix element (as obtained in terms of P-conjugated fields) turns out to amount to an ordinary "CP-conjugated" matrix element. Slightly different is the case of the two (equivalent) neutral weak current couplings entering into (8.4) and (8.5), which look like P-violating terms for the presence of their respective *pseudoscalar* subterms

$$-(g/\cos\theta_W)\sin^2\theta_W(\bar{\psi}_e\gamma^\mu\psi_e)Z_\mu, \qquad +(g/\cos\theta_W)\sin^2\theta_W(\bar{\psi}_{\bar{e}}\gamma^\mu\psi_{\bar{e}})Z_\mu \tag{8.11}$$

(recall that Z_μ is here an axial vector). This is, however, only an apparent P

failure (resulting from the Weinberg mixing): it occurs in the absence of P-breakdown prescriptions and should merely be ascribed to the *mixed* (and not pure pseudoscalar) nature of the "electroweak" effective charge involved in the neutral weak current. In spite of such an asymmetry, the two terms quoted in (8.11) are not prevented from being the P_{in}-conjugated counterparts of each other, just as happens for all the remaining weak coupling terms respectively falling within (8.4) and (8.5): this corresponds to the fact that either term in (8.11) is actually marked by an *identical* (scalar) charge being definite only in magnitude, as can directly be checked by recasting the latter term in the form

$$-(g/\cos\theta_{\mathrm{W}})\,\sin^2\theta_{\mathrm{W}}(\overline{\psi}_{\bar{e}}\gamma^{\mu}\psi_{\bar{e}})(-Z_{\mu})$$

where $(-Z_{\mu})$ is just the P_{in} conjugate of Z_{μ}. More generally, in line with the presence of both scalar and pseudoscalar charges in the model, one has that the overall coupling term (8.3) is left invariant by the *total* covariant charge conjugation CP_{in}, with C acting as in (6.9) and leaving $W_{\mu}, W_{\mu}^{\dagger}$, Z_{μ} unchanged, provided that C (or CP_{in}) is supposed also to apply to A_{μ}, as according to (1.38)

$$C(CP_{\mathrm{in}}):\quad A_{\mu} \to -A_{\mu} \tag{8.12}$$

Hence one consistently obtains that, though A_{μ} is a (true) vector (being left unvaried by P_{in}) and $W_{\mu}, W_{\mu}^{\dagger}$, Z_{μ} are axial vectors (being inverted by P_{in}), the action of CP_{in} upon A_{μ} is the same as upon $W_{\mu}, W_{\mu}^{\dagger}$, Z_{μ}.

It should finally be noticed that either of the equivalent (leptonic and antileptonic) electroweak coupling terms entering into (8.3) allows standard interference between weak and electromagnetic neutral current interactions. This is made possible by the fact that electric charge is still (primarily) a c-number within each *individual* (either leptonic or antileptonic) alternative picture. Actually, the q-number nature of the electric charge "operator" $Q(|q|)$ could fully show up only in the presence of a pseudoscalar charge of the Q^{ch} type, being defined in the Q eigenvector space $\mathcal{S}_{\mathrm{in}}(l_e, \bar{l}_e)$ and directly obeying the anticommutation rule (1.22): in such a case, interference between Q and Q^{ch} dynamics would be strictly forbidden, since each of the two charge operators Q and Q^{ch} necessarily has *null* expectation values pertaining to the single eigenvectors of the other charge operator.

9. CONCLUDING REMARKS

A new approach to the leptonic sector of the electroweak theory[23–25] has been here proposed which is still able to reproduce all the well-established results in that sector,[19] but without need of appealing to those *ad hoc* prescriptions[4–11] normally required to allow for the "maximal P-violation" effect.[2,3]

It relies upon a generalized Dirac field formalism[1] built from introducing a *covariant* charge conjugation—to be identified with proper-mass conjugation[13–16]—according to the Stüeckelberg–Feynman views.[12] The formalism in question is quite consistent with QED and further shows the advantage of giving a natural account of the existence of a chiral phenomenology with no more reason for postulating an actual P failure. In the considered framework, the model of a merely left-handed massless fermion and a merely right-handed related antifermion can be found *spontaneously* and can acquire a *universal* validity (not restricted to the neutrino case) as a result of only *one* pair, and not two independent covariantly charge-conjugated pairs, of (left- and right-handed) chiral field solutions: The massless spin-1/2 fermion and antifermion so obtained, besides naturally agreeing with the neutrino–antineutrino phenomenology, can be theoretically interpreted as two *pseudoscalar*-charge objects being just the *ordinary* mirror image of each other (without any truly "missing" helicity-conjugated counterpart for either of them).

The most immediate consequence of such a retrieval of P mirror symmetry is that an electrically charged fermion and its own antifermion, if taken in their *massless* original version, can at most be thought of with an electric charge that is definite as yet *only in* (squared) *magnitude* and is then conventionally expressible by assigning either single root $\mp|e|$ to *them both*. This statement being allowed for, the weak isospin has been strictly redefined like a quantity whose individual components are ordinary *pseudoscalars* according to (4.2). The peculiar feature of the new weak isospin operator lies in the fact that it can actually generate *two* alternative, $SU(2)_L$ and $SU(2)_R$, symmetry groups, according to *either* γ^5 eigenvalue involved. Within the leptonic sector, such groups in turn give rise to an either (left-handed) lepton or (right-handed) antilepton isodoublet, with (in the former case) antileptons themselves playing the role of (right-handed) $SU(2)_L$ singlets and (in the latter case) leptons themselves playing the role of (left-handed) $SU(2)_R$ singlets. These symmetry groups, along with two appropriate distinct weak hypercharge varieties (Y and $\overline{Y}$) coupled to them, allow the alternate choice of one *suitable* $|e|^2$ root expressing the electric charge (as yet definite only in magnitude) *common* to the (massless) charged lepton–antilepton pair: for instance, the single root $-|e|$ can be assigned to both the charged lepton as a member of a $SU(2)_L$ doublet and the charged antilepton as a $SU(2)_L$ singlet, so that the usual electroweak Gell-Mann–Nishijima formula is identically reproduced (with the charged antilepton in place of the standard *right-handed* charged lepton). Hence two equally allowable (left- and right-handed) covariant pictures—based on the alternative gauge groups $SU(2)_L \otimes U(1)_Y$ and $SU(2)_R \otimes U(1)_{\overline{Y}}$—have been obtained, which can be conveniently merged into a unified $L + R$ formalism.

Another, strictly correlated novelty lies in the fact that the *charged-lepton mass* appearance is now *an essential internal requirement for the generation of the leptonic electromagnetic gauge coupling* instead of being merely an external "accident" connected with a hypothetical coupling to a Higgs boson[26,27]: Here, to become two *actual* eigenstates of electric charge (with opposite eigenvalues), the charged lepton and antilepton *must* indeed be made massive as well. This requirement can (at least formally) be met even *apart from* the Higgs boson existence, just by introducing an effective net *"vacuum" Higgs isodoublet field* merely including the three (massless) Goldstone degrees of freedom and not including the Higgs boson excitation. While being strictly massless and having no quartic self-coupling, the field so introduced may still have a nonzero vacuum expectation value and still cause SSB as well as Weinberg mixing, but with no further effects than mass generation for both the intermediate vector bosons and the charged lepton (antilepton). Of course, such a model cannot fully account yet for the charged-lepton mass spectrum; it can, however, get rid of the questionable standard conjecture of a pure "accidental" origin for that spectrum (as an indirect outcome of the existence of one special coupling constant per each single charged-lepton variety interacting with the Higgs boson).

In the leptonic electroweak scheme proposed, the charged current terms keep their usual chiral-field formal structure, but with a gained theoretical insight into the actual meaning and the general physical legitimacy of a "chiral field" (which may here be soundly introduced even for a massive, and not only for a massless, spin-1/2 point fermion or antifermion). On the grounds of such an insight, those terms can be seen to have a truly *P-invariant* (rather than maximally *P*-violating) form: covariantly speaking, the net *Dirac field V–A* standard currents are now to be viewed as net *chiral field pure-vector* currents. The physical sense of the recovered *P* symmetry is the following: *What, in the light-speed limit, is usually taken as the "missing" P-conjugated counterpart of a left-handed (right-handed) chiral lepton (anti-lepton) may be identified just with the actually existing, right-handed (left-handed) chiral antilepton (lepton)*. This property, which can apply to the chiral behavior of any spin-1/2 point fermion (antifermion), makes the standard weak "*CP*" mirror symmetry substantially reducible to a pure *P* mirror symmetry. The essential fact lies in the *twofold*—either Dirac or chiral—intrinsic nature that is now to be ascribed to a massive spin-1/2 point fermion and antifermion: these, in spite of their customary behavior as scalar-charge eigenstates fully expressed by the *Dirac field*, should strictly be assigned a *pseudoscalar* variety of charges, too, their net looking like pseudoscalar-charge eigenstates fully expressed by the *chiral field*. Such fields can only give a *dual* internal representation of the same fermion and antifermion, as an either *Dirac* or *chiral* particle, respectively. A *Dirac* fermion and the related antifermion individually show a *P*-invariant intrinsic nature: they can be interchanged by scalar-charge conjugation alone. A *chiral* fermion and the related antifermion, on the contrary, are far from showing such a nature: rather, they look just like the ordinary mirror image of each other and can be interchanged by pseudoscalar-charge conjugation alone (already included in *P*). Such antithetic internal attitudes *both*, however, singly respect parity symmetry, even though in two *diametrically opposed* ways,[21] and this fully accounts for the recovered *P*-symmetric form of the charged weak current couplings as well as its conceptual difference from the ordinary *P*-invariant form of the electromagnetic coupling. The same cannot strictly apply to the neutral weak current coupling, where Dirac and chiral field currents interfere (as a result of the Weinberg mixing). Nevertheless, it has been pointed out that such an interference yields only an apparent *P* failure, due to the *mixed* (and not purely pseudoscalar) nature of the "electroweak" effective charge involved.

All this implies no actual deviations from the standard theory at the phenomenological level, except for the *missing* Higgs boson couplings to charged leptons and intermediate vector bosons. It is left to see what further results from extending the new formalism to the (much more intricate) quark sector.

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