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Parity Violation from SUSY Vertex Corrections in Deep Inelastic Scattering

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Abstract

Corrections to the eeγ and qqγ vertices, at the one loop level, from intermediate supersymmetric (SUSY) states are calculated. The induced parity-violating interaction between electrons and nucleons is applied to the calculation of the lowest order SUSY contribution to the asymmetry in deep inelastic e + deuteron scattering in both the 'tree-breaking' (TB) and 'renormalisation-group' (RG) supergravity models. The results indicate that, for a wide range of sparticle masses, the SUSY contributions lie well within constraints implied by present measurements.

1. Introduction

The success of the Glashow–Weinberg–Salam–Ward SU(2)×U(1) unification of electromagnetic and weak interactions has made it fashionable for theoretical physicists to again search for Einstein's holy grail of unified field theories. The next step after the electroweak unification was the introduction of grand unified theories (GUTs) which unify the strong interaction described by the colour SU(3) theory and the SU(2)×U(1) electroweak interactions, and made the spectacular prediction that the proton is unstable. The simplest GUT predicted a proton lifetime of order 10^{31} years, but several experiments have now excluded a lifetime which is this short. However, this theory was remarkably successful in predicting the Weinberg angle which determines the ratio of the weak and electromagnetic interactions and appears as a free parameter in the SU(2)×U(1) theory. This prediction illustrates the hope of those who construct unified theories—that the constraints imposed by unification will enable the prediction of the free parameters of the originally disjoint theories.

There is an aesthetic problem associated with GUTs. The symmetry is realised at energies of order 10^{15} GeV which is the typical mass scale of the theory. When the W mass is calculated it is necessary that the parameters of the theory be 'fine tuned' so that cancellation between loop graphs contributing terms of order 10^{15} GeV produce the electroweak scale of order 100 GeV when the GUT symmetry is broken down to SU(3)×SU(2)×U(1) symmetry of the standard model. This unnatural cancellation in the GUTs is called the hierarchy problem. It was one of the major motivations for introducing SUSY theories in which the gauge group mixes boson and fermion states, and the necessary cancellations now occur naturally because loop contributions of bosons and fermions in the same supermultiplet to the W mass cancel exactly in the SUSY limit. SUSY GUTs have the added advantage of lengthening the predicted

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proton lifetime beyond the present experimental lower bound. [We refer the reader to the recent book by Ross (1985) for a pedagogical review of GUTs and SUSY. Nanopoulous *et al.* (1984) provides a useful review of SUSY theories.]

SUSY GUTs run into their own difficulties unless the supersymmetry is gauged so that it becomes a local symmetry. This has the advantage of introducing gravitational couplings and the resulting theories are referred to as supergravity theories. A useful review of the superspace and superfield formulation of supergravity theories has been given by Gates *et al.* (1983).

A spectacular prediction of SUSY theories is the existence of fermionic partners of the known gauge bosons, and bosonic partners of the known fermions. Thus with the photon, W, Z and gluon there are associated the photino, wino, zino and gluino spin $\frac{1}{2}$ fermions, and with the quarks, leptons and neutrinos there are associated spin zero squarks, sleptons and sneutrinos. The spin-2 graviton is associated with a spin $\frac{3}{2}$ gravitino. In the SUSY limit the particles and their associated sparticles are degenerate in mass, and in the real world the mass splitting of the sparticles from the particles is determined by the SUSY breaking mechanism.

This rich spectrum of new particles has naturally produced an active industry searching for these particles, as well as indirect searches through the corrections the sparticles induce in the standard SU(3)×SU(2)×U(1) phenomenology (see e.g. Godbole 1984; Haber and Kane 1985). At present there is no experimental evidence confirming the realisation of SUSY in Nature. Rather, the null results of existing experiments place constraints on the parameters of the theory.

The search for SUSY takes two broad routes:

(i) direct detection of SUSY particles (e.g. in e^+e^- annihilation);

(ii) precision measurement of second (or higher) order processes to which SUSY is expected to contribute (e.g. $g-2$ factors of μ and e).

In this paper we extend the latter approach by calculating the lowest order SUSY contributions to the parity-violating neutral current effective interaction between leptons and hadrons. The results of this calculation are applied specifically to the observed asymmetry in deep inelastic e+deuteron scattering. Similar work, concerning SUSY contributions to parity violation in nuclei, has been carried out by Suzuki (1982) and extended by Duncan (1983), providing constraints on helicity and isospin mass splittings in the first generation squark spectrum.

Since the measured e+deuteron asymmetry (Prescott *et al.* 1979) is, at best, accurate to about 27%, one expects at best only weak constraints on SUSY parameters. However, the results indicate that useful limits on selectron, squark and sneutrino masses may be obtained from more accurate measurements of parity violation in deep inelastic processes.

2. SUSY Contribution to the Effective Parity-violating Neutral Current Interaction

The particle content of the low energy sectors of minimal $N = 1$ SUSY models gives rise to possible contributions to the parity-violating amplitudes of $ee\gamma$ and $qq\gamma$ vertices at the one loop level. Feynman diagrams for these loop corrections are shown in Fig. 1. As shown, the diagrams separate into two classes: (a) loops containing a single gaugino and (b) those with two intermediate winos. Hence it is necessary to consider only two diagrams in a general manner.

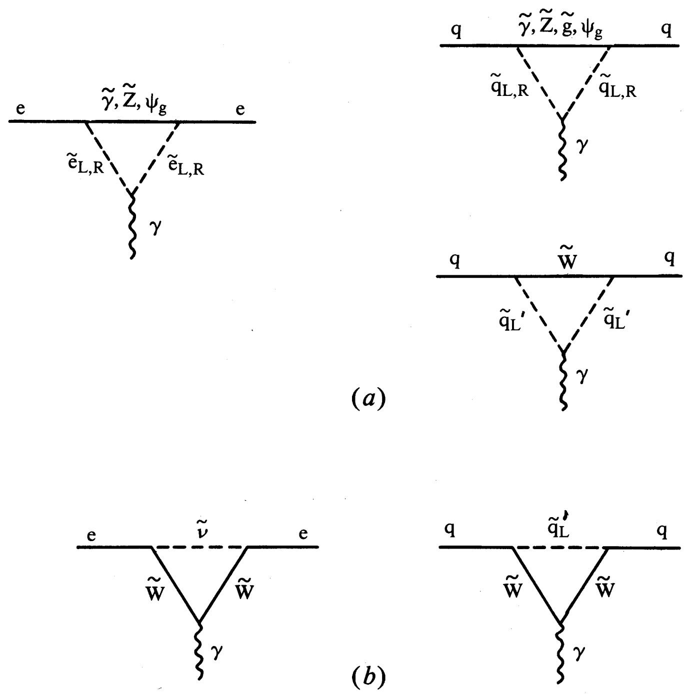

Fig. 1. (*a*) Diagrams modifying the eeγ and qqγ vertices involving a single gaugino; the gauginos $\tilde\gamma$, $\tilde Z$, ψ_g, $\tilde g$ and $\tilde W$ label the photino, zino, gravitino, gluino and wino respectively, whilst $\tilde e_{L,R}$ and $\tilde q_{L,R}$ label the left and right selectrons and squarks. (*b*) Diagrams involving two intermediate winos; the scalar neutrino is labelled by $\tilde\nu$.

In any model, the interaction of a fermion f, its scalar partner $\tilde f$ and a gaugino $\tilde G$ can be written as

$$\mathscr{L}_{f\tilde f\tilde G} = g^{L}(f\tilde f\tilde G)\bar f\tfrac{1}{2}(1+\gamma_5)\tilde G\tilde f_L + g^{R}(f\tilde f\tilde G)\bar f\tfrac{1}{2}(1-\gamma_5)\tilde G\tilde f_R + \text{h.c.}, \tag{1}$$

where $g^L(f\tilde f\tilde G)$ and $g^R(f\tilde f\tilde G)$ are the left and right couplings respectively. The diagrams are straightforward to calculate in terms of (1). We find, working in the limit $m_f^2, q^2 \ll m_{\tilde f}^2$ and the momentum subtraction renormalisation scheme, the parity-violating amplitude for the diagram containing a single gaugino $\tilde G$ to be

$$\Gamma^{\tilde G}_{ff\gamma} = \pm\frac{i}{32\pi^2}\, Q_f e\,|g^{L,R}(f\tilde f\tilde G)|^2 \frac{1}{6m_{\tilde f}^2}\, F_1(\tilde G,\tilde f)\,\bar u_2\gamma_5(q^2\gamma_u - \not{q}q_\mu)u_1 A_\mu\,, \tag{2}$$

and that for the diagram containing two winos

$$\Gamma^{\tilde W}_{ff\gamma} = -\frac{i}{32\pi^2}\, e\,|g^{L}(f\tilde f\tilde W)|^2 \frac{1}{6m_{\tilde f}^2}\, F_2(\tilde W,\tilde f)\,\bar u_2\gamma_5(q^2\gamma_\mu - \not{q}q_\mu)u_1 A_\mu\,, \tag{3}$$

where Q_f is the fermion charge, u_1 and u_2 are spinors representing the initial and final state fermions respectively, q is the momentum transferred to the photon, and the functions F_1 and F_2 are given by

$$F_1(r) = \frac{1}{(1-r)^4}\left(\tfrac{1}{3}-\tfrac{3}{2}r+3r^2-\tfrac{11}{6}r^3+r^3\ln r\right), \tag{4a}$$

$$F_2(r) = \frac{1}{(1-r)^4}\left\{-\tfrac{8}{3}+\tfrac{15}{2}r-6r^2+\tfrac{7}{6}r^3+(3r-2)\ln r\right\}, \tag{4b}$$

and where the notation has been simplified to $F_{1,2}(a, b) = F_{1,2}(r)$, with $r \equiv m_a^2/m_b^2$.

The factor $\pm$ in equation (2) indicates the difference in sign of diagrams containing left (+) and right (−) scalar partners of the fermion f. In general, the helicity states $\tilde{f}_{\mathrm{L,R}}$ will mix to give mass eigenstates $\tilde{f}_{1,2}$ according to

$$\begin{pmatrix}\tilde{f}_1\\ \tilde{f}_2\end{pmatrix} = \begin{pmatrix}\cos\theta_{\mathrm{f}} & \sin\theta_{\mathrm{f}}\\ -\sin\theta_{\mathrm{f}} & \cos\theta_{\mathrm{f}}\end{pmatrix}\begin{pmatrix}\tilde{f}_{\mathrm{L}}\\ \tilde{f}_{\mathrm{R}}\end{pmatrix}, \tag{5}$$

where θ_{f} is the helicity mixing angle. However, in order to reduce the number of undetermined parameters, we have assumed that, for both squarks and sleptons, the mixing is small ($\theta_{\mathrm{f}} \approx 0$) and that the mass eigenstates are approximately $\tilde{f}_{\mathrm{L,R}}$.

Writing the total parity-violating amplitude for the ffγ vertex as

$$\Gamma^{\mathrm{pv}}_{\mathrm{ff}\gamma} = -\frac{\mathrm{i}}{32\pi^2}\,\tfrac{1}{6}F_{\mathrm{ff}\gamma}\,\bar{u}_2\,\gamma_5(\gamma_\mu q^2 - \not{q}\,q_\mu)\,u_1\,A_\mu\,, \tag{6a}$$

we obtain (assuming no generation mixing)

$$F_{\mathrm{ee}\gamma} = e\left(\sum_{i=\mathrm{L,R}}\sum_{\tilde{\mathrm{G}}_{\mathrm{e}}}\frac{\Delta_i}{m^2_{\tilde{\mathrm{e}}_i}}|\mathrm{g}^i(\mathrm{e}\tilde{\mathrm{e}}\tilde{\mathrm{G}}_{\mathrm{e}})|^2\mathrm{F}_1(\tilde{\mathrm{G}}_{\mathrm{e}},\tilde{\mathrm{e}}_i) + \frac{1}{m^2_{\tilde{\nu}}}\sum_j|\mathrm{g}^{\mathrm{L}}(\mathrm{e}\tilde{\nu}\tilde{\mathrm{W}}_j)|^2\mathrm{F}_2(\tilde{\mathrm{W}}_j,\tilde{\nu})\right), \tag{6b}$$

$$\begin{aligned}F_{\mathrm{q}^\alpha\mathrm{q}^\alpha\gamma} = -Q_{\mathrm{q}^\alpha}\,e\Bigg\{&\sum_{i=\mathrm{L,R}}\frac{\Delta_i}{m^2_{\tilde{\mathrm{q}}^\alpha_i}}\left(\sum_{\tilde{\mathrm{G}}_{\mathrm{q}}}|g^i(\mathrm{q}^\alpha\tilde{\mathrm{q}}^\alpha\tilde{\mathrm{G}}_{\mathrm{q}})|^2F_1(\tilde{\mathrm{G}}_{\mathrm{q}},\tilde{\mathrm{q}}^\alpha_i)+4|g^i(\mathrm{q}^\alpha\tilde{\mathrm{q}}^\alpha\tilde{\mathrm{g}})|^2F_1(\tilde{\mathrm{g}},\tilde{\mathrm{q}}^\alpha_i)\right)\\ &+\sum_j\sum_\beta\frac{1-\delta_{\alpha\beta}}{m^2_{\tilde{\mathrm{q}}^\beta_{\mathrm{L}}}}\left(|g^{\mathrm{L}}(\mathrm{q}^\alpha\tilde{\mathrm{q}}^\beta\tilde{\mathrm{W}}_j)|^2\{2F_1(\tilde{\mathrm{W}}_j,\tilde{\mathrm{q}}^\beta_{\mathrm{L}})+F_2(\tilde{\mathrm{W}}_j,\tilde{\mathrm{q}}^\beta_{\mathrm{L}})\}\right)\Bigg\},\end{aligned} \tag{6c}$$

where $\Delta_i \equiv +1, -1$ for $i = \mathrm{L,R}$ respectively, α and β are flavour indices within the generation $\begin{pmatrix}\mathrm{u}\\ \mathrm{d}\end{pmatrix}$, and the gaugino summation sets are

$$\tilde{G}_{\mathrm{e}} = \{\tilde{\gamma}, \tilde{\mathrm{Z}}_k, \psi_{\mathrm{g}}\}\,, \qquad \tilde{G}_{\mathrm{q}} = \{\tilde{\gamma}, \tilde{\mathrm{Z}}_k, \mathrm{W}_j, \psi_{\mathrm{g}}\}\,.$$

Since mass diagonalisation, in most models, results in gaugino mixing within the charged and neutral sectors, we have indicated a summation over all zino and wino states by the indices k and j respectively.

At this stage it is straightforward to write down the contribution from SUSY to the effective parity-violating neutral current interaction (neglecting Z^0 exchange and box diagrams). In terms of the $F_{\mathrm{ff}\gamma}$ (equation 6) we obtain, for a general SUSY model,

$$\mathscr{L}_{\mathrm{pv}}^{(\mathrm{SUSY})} = \frac{1}{32\pi^2}\frac{e}{6}(\tfrac{2}{3}F_{\mathrm{ee}\gamma}\,\bar{e}\gamma_\mu\gamma_5\,e\bar{u}\gamma_\mu\,u - \tfrac{1}{3}F_{\mathrm{ee}\gamma}\,\bar{e}\gamma_\mu\gamma_5\,e\bar{d}\gamma_\mu\,d$$

$$- F_{\mathrm{uu}\gamma}\,\bar{u}\gamma_\mu\gamma_5\,u\bar{e}\gamma_\mu\,e - F_{\mathrm{dd}\gamma}\,\bar{d}\gamma_\mu\gamma_5\,d\bar{e}\gamma_\mu\,e)\,. \tag{7}$$

3. Asymmetry in Deep Inelastic e+deuteron Scattering

The asymmetry in deep inelastic scattering of longitudinally polarised electrons is defined in terms of the left and right differential scattering cross sections $\mathrm{d}\sigma_{\mathrm{L,R}}$ by

$$A \equiv \frac{\mathrm{d}\sigma_{\mathrm{R}} - \mathrm{d}\sigma_{\mathrm{L}}}{\mathrm{d}\sigma_{\mathrm{R}} + \mathrm{d}\sigma_{\mathrm{L}}}\,. \tag{8}$$

If one parametrises the parity-violating neutral current interaction between an electron and the $\binom{\mathrm{u}}{\mathrm{d}}$ generation of quarks by the effective interaction

$$\mathscr{L}_{\mathrm{eff}}^{\mathrm{pv}} = \sum_{q=\mathrm{u,d}} \frac{G_{\mathrm{F}}}{\sqrt{2}}(C_{1q}\,\bar{e}\gamma_\mu\gamma_5\,e\bar{q}\gamma_\mu\,q + C_{2q}\,\bar{e}\gamma_\mu\,e\bar{q}\gamma_\mu\gamma_5 q)\,, \tag{9}$$

then for the scattering from an isoscalar target such as deuterium at $x \geqslant 0{\cdot}2$ (where x is the Bjorken scaling variable), the asymmetry A is given by (Marciano and Sanda 1978)

$$\frac{A(y)}{|q^2|} = a_1 + a_2\,\frac{1-(1-y)^2}{1+(1-y)^2}\,, \tag{10}$$

where the kinematic variable y is given by $y = (E_{\mathrm{i}} - E_{\mathrm{f}})/E_{\mathrm{i}}$ and

$$a_1 = -\frac{3\,G_{\mathrm{F}}}{5\sqrt{2}\,\pi\alpha}(C_{1\mathrm{u}} - \tfrac{1}{2}C_{1\mathrm{d}})\,, \qquad a_2 = -\frac{3\,G_{\mathrm{F}}}{5\sqrt{2}\,\pi\alpha}(C_{2\mathrm{u}} - \tfrac{1}{2}C_{2\mathrm{d}})\,.$$

In obtaining (10) the x dependence has dropped since quark sea effects can be ignored for $x \geqslant 0{\cdot}2$. The asymmetry parameters a_1 and a_2 have been measured at SLAC by Prescott *et al.* (1979), under conditions for which equation (10) is valid. Their model independent fit to data resulted in

$$a_1^{\mathrm{exp}} = (-9{\cdot}7 \pm 2{\cdot}6)\times 10^{-5}\ \mathrm{GeV}^{-2}\,, \qquad a_2^{\mathrm{exp}} = (4{\cdot}9 \pm 8{\cdot}1)\times 10^{-5}\ \mathrm{GeV}^{-2}\,.$$

The first order Weinberg–Salam (WS) model predictions are

$$a_1^{\mathrm{WS}} = -\frac{G_{\mathrm{F}}}{2\pi\sqrt{2}\,\alpha}\,\tfrac{9}{10}(1 - \tfrac{20}{9}\sin^2\theta_{\mathrm{W}}) = -8{\cdot}26\times 10^{-5}\ \mathrm{GeV}^{-2}\,,$$

$$a_2^{\mathrm{WS}} = -\frac{G_{\mathrm{F}}}{2\pi\sqrt{2}\,\alpha}\,\tfrac{9}{10}(1 - 4\sin^2\theta_{\mathrm{W}}) = -1{\cdot}95\times 10^{-5}\ \mathrm{GeV}^{-2}\,,$$

for $\sin^2\theta_{\mathrm{W}} = 0{\cdot}22$. The standard model is consistent with the measured asymmetry to within the uncertainties quoted. Although the experimental errors are large (especially for a_2), one can still give meaningful constraints on the contributions to a_1 and a_2 from the SUSY neutral current interaction (equation 7) in terms of the WS predictions (we ignore the second order WS contributions, and the bizarre possibility that a large SUSY contribution is cancelled by another large anomalous contribution); for example

$$-4\cdot0\times10^{-5}\ \mathrm{GeV}^{-2} \lesssim a_1^{\mathrm{SUSY}} \lesssim 1\cdot2\times10^{-5}\ \mathrm{GeV}^{-2}. \tag{11}$$

In the case of a_2 the sign and magnitude of the measured value is unclear, but we can safely say that

$$|a_2^{\mathrm{SUSY}}| \lesssim 1\times10^{-4}\ \mathrm{GeV}^{-2}. \tag{12}$$

Using (7) and the above numerical constraints, we obtain

$$-4\cdot0\times10^{-5}\ \mathrm{GeV}^{-2} \lesssim -\frac{e}{32\pi^2}\frac{1}{12\pi\alpha}F_{\mathrm{ee}\gamma} \lesssim 1\cdot2\times10^{-5}\ \mathrm{GeV}^{-2}. \tag{13a}$$

$$\left|\frac{e}{32\pi^2}\frac{1}{10\pi\alpha}(F_{\mathrm{uu}\gamma}-\tfrac{1}{2}F_{\mathrm{dd}\gamma})\right| \lesssim 1\times10^{-4}\ \mathrm{GeV}^{-2}. \tag{13b}$$

The constraints (13) are completely general with respect to SUSY models. Therefore, more accurate measurements of the parameters a_1 and a_2 may make it possible to illustrate differences between models whilst deriving parameter space constraints.

4. Low Energy Phenomenology of Supergravity Models

For definiteness we consider two classes of supergravity GUT models based on the gauge group SU(5) (Chamseddine *et al.* 1982; Alvarez-Gaumé *et al.* 1983). In these models the spontaneous breaking of SU(2)×U(1) is a residual effect caused by the presence of soft terms in the low energy effective potential induced by couplings to the $N = 1$ supergravity. Possible mechanisms for the SU(2)×U(1) breakdown are:

(i) breaking at the tree level (TB model);

(ii) breaking through renormalisation group loop corrections (RG model).

Spontaneous breaking of SU(2)×U(1) in TB models is achieved by introducing an extra singlet superfield $\hat{U}$ (where the circumflex denotes a superfield). In the RG model there are a number of ways in which large corrections can be generated to break the SU(2)×U(1) symmetry. We concentrate on the case where a large top quark mass is assumed, i.e. $m_{\mathrm{t}} \approx$ 100–200 GeV (Alvarez-Gaumé *et al.* 1983). For this range of m_{t} the Higgs mass m_{H}^2 goes negative in the renormalisation group equations at a SU(2)×U(1) breaking scale of the order of 100 GeV. We note that the top quark has not yet been unambiguously identified and our assumption regarding its mass is not in conflict with experiment.

Following Nath *et al.* (1984) we assume a common Higgs structure of one pair of superfield doublets $\hat{H}^i$ and $\hat{H}'_i$ (i.e. the simplest case) which thus enables a simultaneous parametrisation of the TB and RG models. Although, strictly speaking, one needs at least two pairs of Higgs doublets in the TB model to counteract loop instabilities introduced by the singlet $\hat{U}$, two of the Higgs particles become heavy and the low energy Higgs structure is essentially that of the RG model with one pair of doublets and a singlet $\hat{U}$.

The parametrisation by Nath *et al.* (1984), which provides the useful interpolation between the two models, begins by defining the Higgs angle θ_H by

$$\tan\theta_H \equiv \langle H_2'\rangle/\langle H^2\rangle.$$

The superpotential contains couplings $\hat{H}^i\hat{H}_i'$ and $\hat{U}\hat{H}^i\hat{H}_i'$ whose strengths are characterised by μ and μ' respectively. Table 1 summarises the differences between the two classes of supergravity models in terms of these parameters.

Table 1. Typical parameter values for TB and RG models

m_g is the gravitino mass

	TB	RG
θ_H	40–50°	10–20°
μ	$\sim m_g$	$\sim\theta_H\, m_g$
μ'	$\sim m_g$	0

Diagonalisation of the charged gaugino and higgsino fields leads to two charged wino states $\tilde{W}_{(+)}$ and $\tilde{W}_{(-)}$ with masses $m_{\tilde{W}(\pm)}$ lying either side of the W boson and given by (Nath *et al.* 1984)

$$m_{\tilde{W}(\pm)} = \tfrac{1}{2}[\{4\nu_+^2+(\mu-1\cdot7\,m_{\tilde{\gamma}})^2\}^{\frac{1}{2}} \pm\{4\nu_-^2+(\mu+1\cdot7\,m_{\tilde{\gamma}})^2\}^{\frac{1}{2}}], \qquad (14)$$

where $\nu_\pm = \surd\frac{1}{2}\, M_W(\cos\theta_H\pm\sin\theta_H)$ and $m_{\tilde{\gamma}}$ is the photino mass.

The neutral gaugino mass matrix can be diagonalised analytically if the photino is an approximate mass eigenstate and $(\mu\cos2\theta_H)^2 \ll m_Z^2$ (which holds for both RG and TB models). One finds, in addition to the photino, four zino mass eigenstates: the states $\tilde{Z}_{(\pm)}$ lying either side of the Z^0 boson with masses (Nath *et al.* 1984)

$$m_{\tilde{Z}(\pm)} = \{m_Z^2+\tfrac{1}{4}(\mu\sin2\theta_H-1\cdot5\,m_{\tilde{\gamma}})^2\}^{\frac{1}{2}}\pm\tfrac{1}{2}(\mu\sin2\theta_H+1\cdot5\,m_{\tilde{\gamma}}), \qquad (15)$$

and $\tilde{Z}_{(3,4)}$ with masses* (Nath *et al.* 1984)

$$m_{\tilde{Z}(3,4)} = (\tfrac{1}{4}\mu^2\sin^2 2\theta_H+\mu'^2)^{\frac{1}{2}}\pm\tfrac{1}{2}\mu\sin2\theta_H\,. \qquad (16)$$

The existence of $\tilde{Z}_{(4)}$ is due entirely to the extra singlet field $\hat{U}$ in the TB model (from the values of μ and μ' in Table 1 one sees that $\tilde{Z}_{(3,4)}$ are usually quite heavy in this model). In the RG model one has the states $\tilde{Z}_{(\pm)}$ and the so-called 'twilight zino' $\tilde{Z}_{(3)}$, which is generally light and couples weakly.

The masses of the first generation of squarks and sleptons are given by (Nath *et al.* 1984)

$$m^2_{\tilde{f}(L,R)} = m_g^2\mp(I_w^3-Q_f\sin^2\theta_W)m_Z^2\cos2\theta_H\,, \qquad (17)$$

* Strictly speaking, one should include a parameter μ'' corresponding to the cubic self-interaction of the singlet $\hat{U}$ in the superpotential of the TB model, which manifests itself in the masses $m_{\tilde{Z}(3,4)}$ and the couplings of the four zino states. However, the effects of this extra parameter on the $\tilde{Z}_{(\pm)}$ couplings are very small and, since the $\tilde{Z}_{(3,4)}$ couple extremely weakly, we neglect μ'' in our calculations.

where I^3_w is the third component of the weak isospin, Q_f is the charge of the fermion f, and m_g is the gravitino mass.

The couplings $|g^{L,R}(f\tilde{f}\tilde{G})|^2$, as defined by equation (1), are given in Table 2. The zino and wino couplings involve mixing factors arising from the diagonalisation matrices given by

$$\Lambda_k = m_Z \frac{\tilde{\lambda}_k}{D_k}, \qquad \tilde{\lambda}_k = \lambda_k + \mu \sin 2\theta_H - \mu'^2/\lambda_k,$$

$$D_k = [m_Z^2 \tilde{\lambda}_k^2 + (\lambda_k - 1\cdot 5 m_{\tilde{\gamma}})^2 \{\tilde{\lambda}_k^2 + \mu^2 \cos^2 2\theta_H (1 + \mu'^2/\lambda_k^2)\}]^{\frac{1}{2}},$$

$$\lambda_\pm = \pm m_{\tilde{Z}(3,4)}, \qquad \lambda_{3,4} = \mp m_{\tilde{Z}(3,4)},$$

$$\gamma_\pm = \beta_+ \pm \beta_-, \qquad \tan\beta_\pm = \frac{1}{2\nu_\pm}(\mu \mp 1\cdot 7 m_{\tilde{\gamma}}).$$

The gluino coupling is fixed by the strong QCD coupling constant g_s evaluated at the loop scale.

Table 2. Relevant couplings (as defined in equation 1)

$f\tilde{f}\tilde{G}$	$\lvert g^L(f\tilde{f}\tilde{G})\rvert^2$	$\lvert g^R(f\tilde{f}\tilde{G})\rvert^2$
$u\tilde{d}\tilde{W}_{(\pm)}$	$\frac{e^2}{\sin^2\theta_W}\begin{pmatrix}\cos^2\gamma_+\\ \sin^2\gamma_+\end{pmatrix}$	
$d\tilde{u}\tilde{W}_{(\pm)}$	$\frac{e^2}{\sin^2\theta_W}\begin{pmatrix}\cos^2\gamma_-\\ \sin^2\gamma_-\end{pmatrix}$	
$e\tilde{\nu}\tilde{W}_{(\pm)}$	$\frac{e^2}{\sin^2\theta_W}\begin{pmatrix}\sin^2\gamma_-\\ \cos^2\gamma_-\end{pmatrix}$	
$f\tilde{f}\tilde{\gamma}$	$2e^2 Q_f^2$	$2e^2 Q_f^2$
$f\tilde{f}\tilde{Z}_k$	$\frac{2e^2(I^3_w - Q_f \sin^2\theta_W)^2 \Lambda_k^2}{\sin^2\theta_W \cos^2\theta_W}$	$\frac{2e^2(I^3_w - Q_f \sin^2\theta_W)^2 \Lambda_k^2}{\sin^2\theta_W \cos^2\theta_W}$
$q\tilde{q}\tilde{g}$	$2g_s^2$	$2g_s^2$

5. Results and Discussion

The asymmetry parameters $a_1^{\rm SUSY}$ and $a_2^{\rm SUSY}$ are plotted against the gravitino mass in Fig. 2. These curves are only illustrative in the sense that they will vary slightly with the intrinsic model parameters θ_H, μ and μ'. However, these parameters are constrained with each model, as indicated in Table 1, at least to the extent that the curves will remain the same order of magnitude shown. As the gravitino is almost non-interacting in a local SUSY theory we have ignored its contributions.*

* This may not be the case for a spontaneously broken global SUSY theory in which the gravitino may be light and couple with strength $|g^{L,R}(f\tilde{f}\psi_g)|^2 \approx 2(m^2_{\tilde{f}(L,R)}/d)^2$, where d measures the scale of SUSY breaking. However, using a lower bound of $d \gtrsim 1350$ GeV^2, obtained from $(g-2)_\mu$ SUSY constraints (Barbieri and Mainani 1982), we find that the gravitino contributions to $a_1^{\rm SUSY}$ and $a_2^{\rm SUSY}$ are of the order $(10^{-3})\times 10^{-5}$ GeV^{-2}.

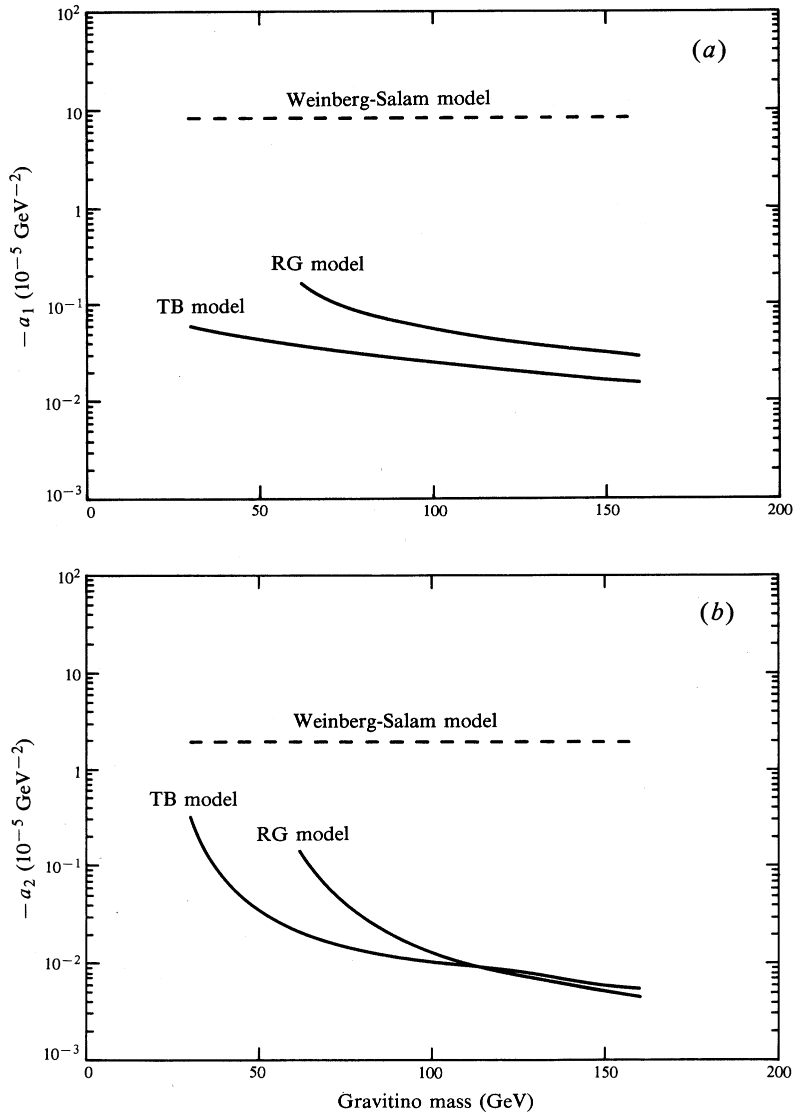

Fig. 2. Corrections to the parity-violating asymmetry parameters (*a*) a_1 and (*b*) a_2 in the TB and RG models (for $m_{\tilde{\gamma}} = 0$ and $m_{\tilde{g}} = 10$ GeV).

In the TB model the helicity splittings of the squarks and selectrons are maximised by choosing $\theta_H(\mathrm{TB}) = 40°$ (or $\theta_H = 50°$) and are degenerate with m_g for $\theta_H = 45°$. The value of θ_H is not so critical in the RG model and we have taken a mid range value of $\theta_H(\mathrm{RG}) = 15°$. Naive lower bounds on the gravitino mass are obtained from the condition $m_{\tilde{\nu}} \geqslant 0$, in which case

$$m_g \geqslant 28 \text{ GeV} \quad (\text{TB with } \theta_H = 40°)$$

$$\geqslant 62 \text{ GeV} \quad (\text{RG with } \theta_H = 15°)\,.$$

In reality, one expects the gravitino mass to be ~ 100 GeV in both models [in fact, for the RG model, a lower bound on the gravitino mass of about $O(m_W)$ can be obtained for a range of top quark masses and GUT parameters].

Preserving the neutral gaugino mass diagonalisation condition, $(\mu \cos 2\theta_H)^2 \ll m_Z^2$, for the values of θ_H(TB) and θ_H(RG) chosen requires

$$\mu(\mathrm{TB}) \ll 540\ \mathrm{GeV}, \qquad \mu(\mathrm{RG}) \ll 100\ \mathrm{GeV}. \tag{18}$$

Hence, we have used the 'maximal' values

$$\mu(\mathrm{TB}) = 100\ \mathrm{GeV}, \qquad \mu(\mathrm{RG}) = 10\ \mathrm{GeV}. \tag{19}$$

Variations of order 50 GeV on μ(TB) and 10 GeV on μ(RG) did not significantly alter the results.

The parameter μ' in the TB model is expected to be of the same order as μ and so we have taken $\mu' = 100$ GeV. However, since the $\tilde{Z}_{(3)}$ and $\tilde{Z}_{(4)}$ couplings are small, the results are largely independent of the value of μ'.

The photino is expected to be by far the lightest gaugino, as we have calculated for $m_{\tilde{\gamma}} = 0$. The effect of increasing $m_{\tilde{\gamma}}$ is to increase the curves slightly; however, for $m_{\tilde{\gamma}} \lesssim 20$ GeV, there is little alteration in the general order of magnitude appearance of the results. Since increasing the gluino mass tends to decrease the magnitude of the curves we have selected $m_g = 10$ GeV, a reasonable value for this parameter [current estimates from $p\bar{p}$ results at CERN and cosmological arguments favour $m_{\tilde{g}}$, $m_{\tilde{q}} \geqslant O(40\text{–}50\ \mathrm{GeV})$ (Ellis 1985), however, $m_{\tilde{g}} \sim 10$ GeV and $m_{\tilde{q}} \sim 100$ GeV cannot as yet be ruled out (Barnett *et al.* 1985)].

As can be seen from Fig. 2, the SUSY contributions to a_1 and a_2 lie well within the constraints (11) and (12) for a wide range of model parameters in TB and RG supergravity GUT models. It appears then that at the present level of experimental accuracy, signatures of SUSY from parity violation in deep inelastic scattering are well hidden (being essentially a second order effect). In terms of the two classes of supergravity models considered here, the measurement of asymmetry in $e+d \rightarrow e+x$ would have to be improved by about two orders of magnitude before signals of SUSY would become evident.

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