τ polarization asymmetry in $B \to X_s \tau^+ \tau^-$ in SUSY models with large $\tan\beta$

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Rare B decays provide an opportunity to probe for new physics beyond the standard model. The effective Hamiltonian for the decay $b \to s l^+ l^-$ predicts the characteristic polarization for the final state lepton. Lepton polarization has, in addition to a longitudinal component P_L, two orthogonal components P_T and P_N lying in and perpendicular to the decay plane. In this article we perform a study of the τ-polarization asymmetry in the case of SUSY models with large $\tan\beta$ in the inclusive decay $B \to X_s \tau^+ \tau^-$. [S0556-2821(99)06221-9]

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Recent progress in experiment and theory has made flavor changing neutral current (FCNC) B decays a stringent test of the standard model (SM) and a powerful probe of physics beyond the standard model. The first observations [1] of the inclusive and exclusive radiative decays $B \to X_s \gamma$ and $B \to K^* \gamma$ have placed the study of rare B decays on a new footing. The observation of $b \to s\gamma$ by CLEO puts very strong constraints on various new physics beyond the standard model. In the case of $B \to X_s \gamma$ CLEO Observations give very strong constraints on the charged Higgs boson mass in the two Higgs doublet model. But in minimal supersymmetric standard model (MSSM) these constraints becomes a bit relaxed because of various cancellations between different superparticle contributions. It is therefore important to study the sensitivity of other FCNC processes to SUSY.

Recently the inclusive decay of $B \to X_s l^+ l^-$ [2,3] received considerable attention as a testing ground of SM and new physics. The experimental situations of these decays is very promising with e^+e^- and hadronic colliders closing on the observation of exclusive models with $l = \mu$ and e final states, respectively. In this decay we can observe various kinematical distributions associated with a final state lepton pair such as lepton pair invariant mass spectrum, lepton pair forward backward asymmetry, etc. Recently another observable, τ polarization asymmetry, for the $B \to X_s \tau^+ \tau^-$ mode has also been proposed by Hewett [4] which can again be used for more strict checking of effective Hamiltonian governing the decay. In another work [5] attention has been drawn to the fact that apart from longitudinal polarization of lepton there can be two other orthogonal components of polarizations which are proportional to m_l/m_b and hence are important for τ. These components of polarizations, namely, the component in the decay plane (P_T, transverse polarization) and the component normal to decay plane (P_N, normal polarization).[1] In this paper we will try to examine the sensitivity of these observables with respect to new physics, i.e., MSSM.

Among models for physics beyond standard model supersymmetry (SUSY) is considered to be the most promising candidate. The minimal extension of the standard model (MSSM) involves chiral superfields Q, U^c, D^c, L, E^c, H_1, and H_2 which transforms under $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ as

$$Q \equiv (3,2,1/2), \quad U^c \equiv (\bar{3},1,-2/3),$$
$$D^c \equiv (\bar{3},1,1/3), \quad L \equiv (1,2,-1/2),$$
$$E^c \equiv (1,1,1), \quad H_1 \equiv (1,2,-1/2),$$
$$H_2 \equiv (1,2,1/2). \tag{1}$$

The superpotential in MSSM in terms of these superfields are

$$W = h_U^{ij} Q_i U_j^c H_2 + h_D^{ij} Q_i D_j^c H_1 + \mu H_1 H_2 + h_e^{ij} L_i E_j^c H_1, \tag{2}$$

where i,j denote generation indices ($i,j = 1,2,3$), and μ and h's are parameters of MSSM. Supersymmetry is broken softly in MSSM. At a large grand-unified scale M_G the bilinear terms have the structure

$$M_{\text{soft}}^{(2)} = \sum_i m_i^2 |y_i|^2 + \frac{1}{2} \sum_j (M_j \lambda_i \lambda_j + \text{H.c.}), \tag{3}$$

where y_i's are the scalar components of the chiral superfields and $\lambda_1, \lambda_2, \lambda_3$ are the two component gaugino fields of $\mathrm{U}(1)_Y$, $\mathrm{SU}(2)_L$, and $\mathrm{SU}(3)_c$, $m_i, M_i' s$ are parameters. The trilinear soft breaking term is

$$M_{\text{soft}}^{(3)} = mA[h_U Q^s U^s H_2^s + h_D Q^s D^s H_1^s + h_E L^s E^s H_1^s] + Bm\mu H_1^s H_2^s + \text{H.c.}, \tag{4}$$

where the superscript s denote the scalar component of the corresponding superfield and the generation index is suppressed in Eq. (4). A and B are constants and m is a scale factor. At scale $\sim M_W$, the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ symmetry is broken spontaneously by the H_1, H_2 developing a nonzero vacuum expectation value.

$$\langle H_1 \rangle = \begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \quad \langle H_2 \rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix}. \tag{5}$$

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[1]Different combinations of the Wilson coefficients describing the decay and are thus useful for comparing theory with experimental data.

The quantity $\tan\beta = v_2/v_1$, thus enters as another parameter in MSSM.

The MSSM in its general form has far too many parameters for it to be used in phenomenology in any meaningful way. Most applications have considered MSSM in the context of minimal spontaneously broken $N=1$ supergravity (SUGRA). This implies that at the Planck scale all the scalar masses have an universal value ($m_i = m$) as do the gauginos ($M_i = M$). At M_G we thus have five parameters (apart from gauge and Yukawa couplings and $\tan\beta$) A, B, μ, m, and M. Using renormalization group equations these parameters can be scaled down to the scale M_W. The condition that at scale M_W, the $SU(2)_L \times U(1)_Y$ symmetry breaks down to $U(1)_{em}$, via the spontaneous symmetry breaking (SSB) condition Eq. (5), reduces the number of independent parameters 2. However, as discussed in Ref. [6], we use a more relaxed SUGRA model which requires the degeneracy of soft SUSY-breaking mass in the scalar squark sector and separately in the Higgs boson sector; thus in Eq. (3) $m_i = m_0$ for squarks and $m_i = \Delta_0$ for the Higgs boson. This, as has been discussed by Goto *et al.* [6], is sufficient to ensure an important constraint, namely adequate suppression of $K^0 - \overline{K}^0$ mixing.

The MSSM has been used to study various rare decays such as $b \to s l^+ l^-, b \to s \nu \bar{\nu}, K^0 \to \pi^0 l^+ l^-$ using the known results of $b \to s\gamma$ [1] as a constraint on the parameter space [7,8]. It was also observed that very large value of $\tan\beta$ is still allowed [7–10] It has been pointed out recently by [7,11] that for large $\tan\beta$, which is allowed by the constraining condition, the process $b \to s l^+ l^-$ can also proceed via exchange of neutral Higgs bosons (NHB) h^0, H^0 and A^0. These exchanges lead to additional amplitudes which scale like $m_b m_l \tan^3\beta$ and this can give considerable enhancement of processes such as $B_s \to \mu^+\mu^-, B \to X_s l^+ l^-$, etc. [7,11]. For $l = \tau$, these NHB contributions for large $\tan\beta$ will be even more significant. In this paper we will try to estimate τ-polarization parameters including NHBs contributions.

We start by writing down the QCD improved effective Hamiltonian for the process $B \to X_s l^+ l^-$ [11]:

$$\mathcal{H} = \frac{\alpha G_F}{\sqrt{2}\pi} V_{tb} V_{ts}^* \Big[C_9^{eff} (\bar{s}\gamma_\mu P_L b)\bar{l}\gamma^\mu l + C_{10}(\bar{s}\gamma_\mu P_L b) \times \bar{l}\gamma^\mu\gamma^5 l - 2 C_7^{eff}\, \bar{s} i \sigma_{\mu\nu} \frac{p^\nu}{p^2}(m_b P_R + m_s P_L) b \bar{l}\gamma^\mu l + C_{Q_1}(\bar{s}P_R b)\bar{l}l + C_{Q_2}(\bar{s}P_R b)\bar{l}\gamma_5 l \Big] \quad (6)$$

with $P_{L,R} = \frac{1}{2}(1 \mp \gamma_5), p = p_+ + p_-$ and where the sum of the momentum of l^+ and l^-, C_9^{eff}, C_{10}, and C_7^{eff} are Wilson coefficients given in Refs. [5,12]. C_{Q_1} and C_{Q_2} are new Wilson coefficients which are absent in the standard model but arises in MSSM due to NHB exchange. Their values are given in Ref. [11].The C's all receive contributions from diagrams involving SUSY particles. However, as has been pointed out in Refs. [6,8], the various SUSY contributions to C_7, C_9, and C_{10} have large cancellations amongst themselves leading to only mild changes in their values relative to SM. C_{Q_1} and C_{Q_2}'s ,which have only SUSY contributions, for certain regions of allowed parameter space (space allowed by $b \to s\gamma$) can be comparable to magnitude of C_{10}. We, however, include the SUSY contributions to all Wilson coefficients as given in Refs. [6,8] for our numerical estimates.

$B \to X_s l^+ l^-$ also receives large long distance contributions from tree level process associated with $c\bar{c}$ resonances in intermediate states, i.e., with chain reaction $B \to X_s + \Psi \to X_s l^+ l^-$. These resonant contributions can be incorporated into lepton pair invariant mass spectrum according to prescription of Ref. [13] by employing Breit-Wigner form of the resonance propagator. This produces an additional contribution to C_9^{eff} of the form

$$\frac{-3\pi}{\alpha^2} \sum_{V = J/\psi, \psi', \ldots} \frac{M_V Br(V \to l^+ l^-) \Gamma_{\text{total}}^V}{(s - M_V^2) + i\Gamma_{\text{total}}^V M_V}, \quad (7)$$

where the properties of the vector mesons are given in a table in Ref. [5]. There are six known resonances in the $c\bar{c}$ system that can contribute to the decay modes $B \to X_s l^+ l^-$. Of these, all except the lowest $J/\psi(3097)$ contribute to the channel $B \to X_s \tau^+\tau^-$, for which the invariant mass of lepton pair is $s > 4m_\tau^2$, i.e., greater then τ pair production threshold.[2]

The differential decay rate for $B \to X_s \tau^+\tau^-$ is then

$$\frac{dB(B \to X_s \tau^+\tau^-)}{d\hat{s}} = \frac{G_F^2 m_b^5}{192\pi^3} \frac{\alpha^2}{4\pi^2} |V_{tb} V_{ts}^*|^2 \lambda^{1/2}(1, \hat{s}, \hat{m}_s^2) \sqrt{1 - \frac{4\hat{m}_l^2}{\hat{s}}} \triangle, \quad (8)$$

where factors λ and $\triangle$ are defined by

$$\lambda(a,b,c) = a^2 + b^2 + c^2 - 2(ab + bc + ac) \quad (9)$$

and

[2]As given in references the prescription Eq. (7) for the resonance contribution implies an inclusive direct J/ψ production rate $Br(B \to J/\psi X_s) = 0.15$ that is ~ 5 times smaller than the measured J/ψ rate. This is corrected by the introduction of a phenomenological factor of $\kappa_v \approx 2$ multiplying the Breit-Wigner function in Eq. (7). For our results we use $\kappa_v = 2.35$.

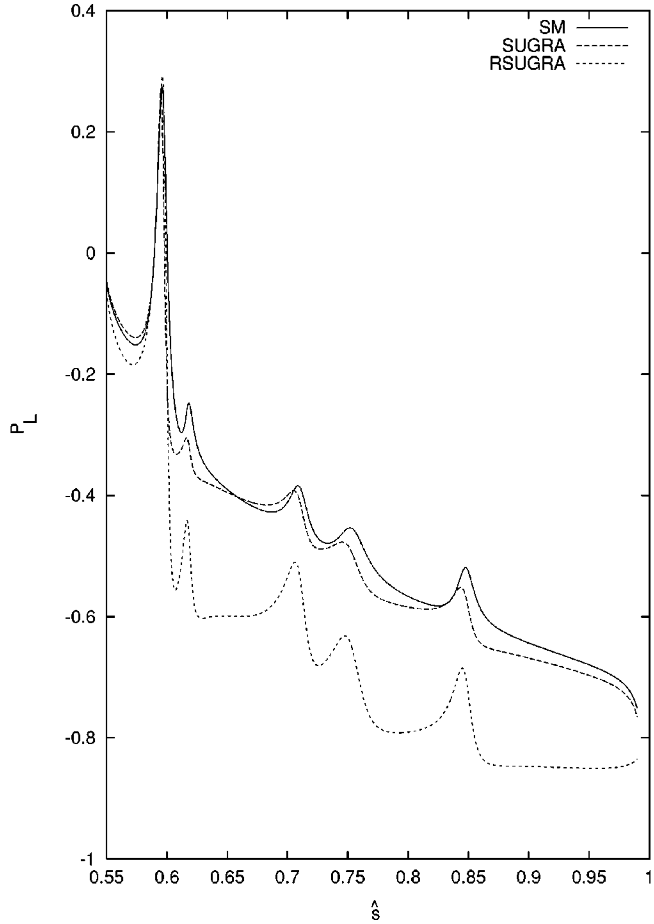

FIG. 1. Longitudinal polariztion asymmetry with $\hat{s}$. The parameters taken are $\tan\beta = 30; m_0 = M = 130; A = -1$, for the relaxed SUGRA (RSUGRA) model $m_A = 120$. All masses are in GeV.

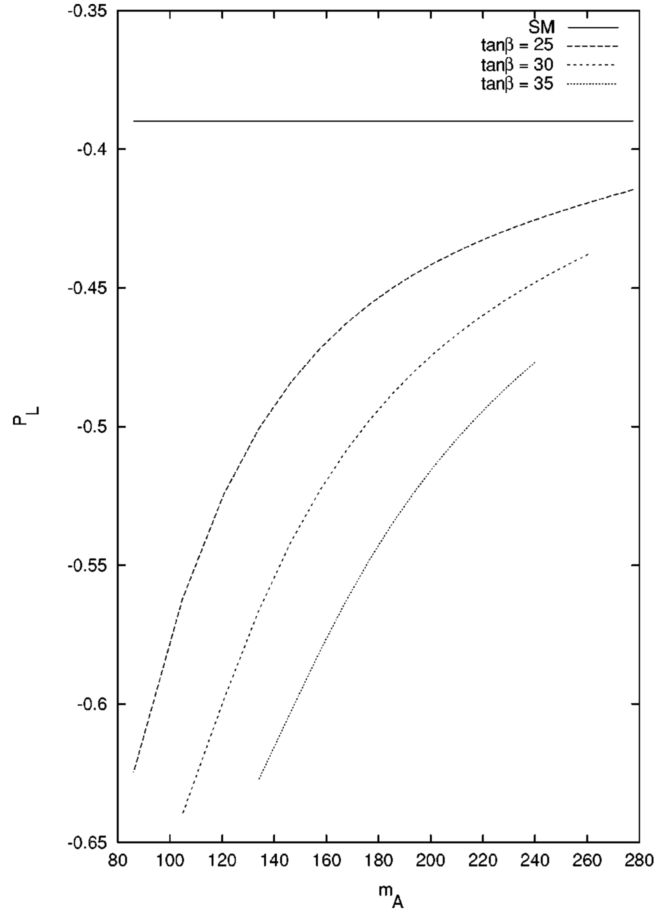

FIG. 2. Longitudinal polarization asymmetry with m_A in relaxed SUGRA model. Other parameters are $m_0 = M = 130; A = -1; \hat{s} = 0.65$. All masses are in GeV.

$$\triangle = \left\{ \left(\frac{4}{\hat{s}} |C_7^{\rm eff}|^2 F_1(\hat{s},\hat{m}_s^2) + 12 {\rm Re}(C_7^{\rm eff} C_9^{\rm eff}) F_2(\hat{s},\hat{m}_s^2) \right) \times \left(1 + \frac{2\hat{m}_l^2}{\hat{s}} \right) + |C_9^{\rm eff}|^2 F_3(\hat{s},\hat{m}_s^2,\hat{m}_l^2) + |C_{10}|^2 F_4(\hat{s},\hat{m}_s^2,\hat{m}_l^2) + \frac{3}{2} |C_{Q_1}|^2 F_5(\hat{s},\hat{m}_s^2)(\hat{s} - 4\hat{m}_l^2) + \frac{3}{2} |C_{Q_2}|^2 F_6(\hat{s},\hat{m}_s^2) + 6 C_{Q_2} C_{10} \hat{m}_l F_7(\hat{s},\hat{m}_s^2) \right\}, \tag{10}$$

with

$$F_1(\hat{s},\hat{m}_s^2) = 2(1+\hat{m}_s^2)(1-\hat{m}_s^2)^2 - \hat{s}(1+14\hat{m}_s^2+\hat{m}_s^4) - \hat{s}^2(1+\hat{m}_s^2), \tag{11}$$

$$F_2(\hat{s},\hat{m}_s^2) = (1-\hat{m}_s^2)^2 - \hat{s}(1+\hat{m}_s^2), \tag{12}$$

$$F_3(\hat{s},\hat{m}_s^2,\hat{m}_l^2) = (1-\hat{m}_s^2)^2 + \hat{s}(1-\hat{m}_s^2) - 2\hat{s}^2 + \frac{2\hat{m}_l^2}{\hat{s}} \times [(1+\hat{m}_s^2)\hat{s} + (1-\hat{m}_s^2)^2 - 2\hat{s}], \tag{13}$$

$$F_4(\hat{s},\hat{m}_s^2,\hat{m}_l^2) = (1-\hat{m}_s^2)^2 + \hat{s}(1-\hat{m}_s^2) - 2\hat{s}^2 + \frac{2\hat{m}_l^2}{\hat{s}} \times [-5(1+\hat{m}_s^2)\hat{s} + (1-\hat{m}_s^2)^2 + 4\hat{s}], \tag{14}$$

$$F_5(\hat{s},\hat{m}_s^2) = 1 + \hat{m}_s^2 - \hat{s}, \tag{15}$$

$$F_6(\hat{s},\hat{m}_s^2) = \hat{s}(1 + \hat{m}_s^2 - \hat{s}), \tag{16}$$

$$F_7(\hat{s},\hat{m}_s^2) = 1 - \hat{m}_s^2 - \hat{s}, \tag{17}$$

where we have used the notion that $\hat{s} = p^2/m_b^2, \hat{m}_i = m_i/m_b$. This matches the result of Ref. [11].

Now we discuss the final state lepton polarization. The polarized cross sections are obtained by introducing the spin projection operator. For l^-,

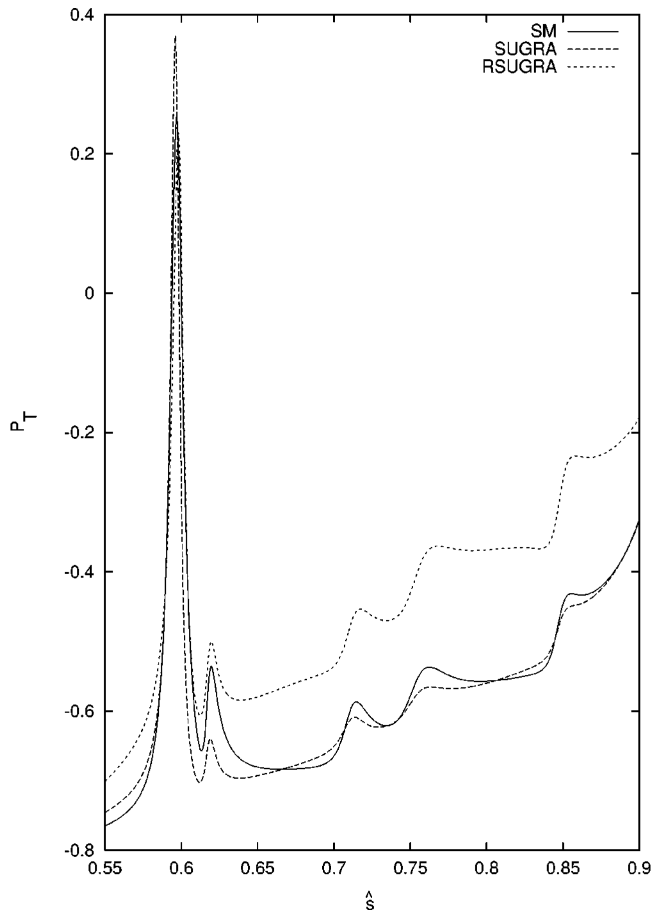

FIG. 3. Transverse polarization asymmetry with $\hat{s}$. Other parameters are $\tan\beta=30; m_0=M=130; A=-1$, for the relaxed SUGRA (RSUGRA) model $m_A=120$. All masses are in GeV.

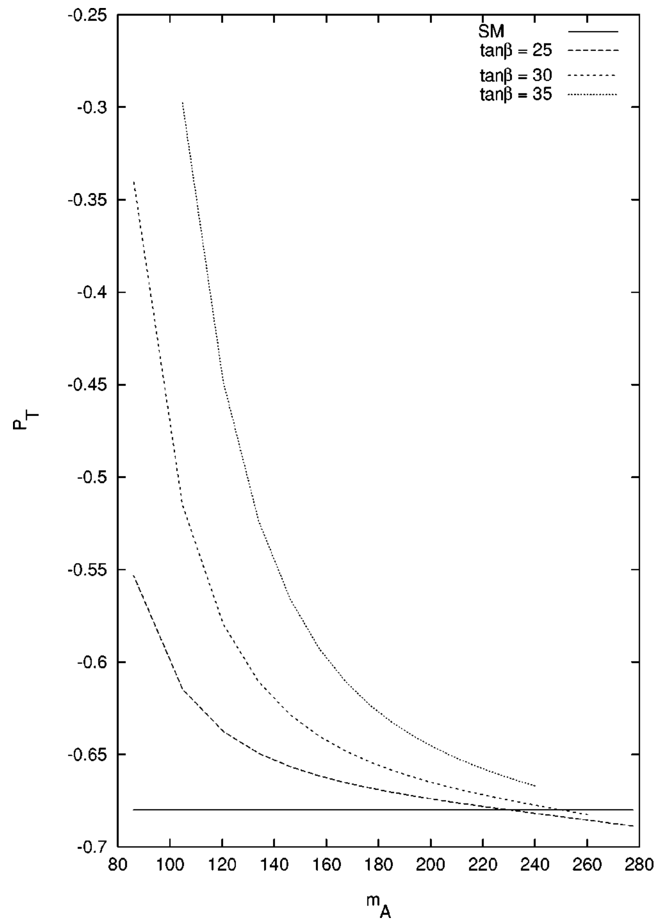

FIG. 4. Transverse polarization asymmetry with m_A in relaxed SUGRA model. Other parameters are $m_0=M=130; A=-1; \hat{s}=0.65$. All masses are in GeV.

$$P=\frac{1}{2}(1+\gamma_5 \not{N}_i), \quad i=L,T,N. \tag{18}$$

$(N_\mu)_i$ here are four vectors satisfying $Np_-=0$ and $N^2=-1$. In general,

$$(N_\mu)_L=\left(\frac{|\hat{p}_-|}{m_l}, \frac{p_-^0}{m_l}\hat{e}_L\right), \tag{19}$$

$$(N_\mu)_T=(0,\hat{e}_T), \tag{20}$$

$$(N_\mu)_N=(0,\hat{e}_N) \tag{21}$$

where

$$\hat{e}_L=\hat{p}_-, \tag{22}$$

$$\hat{e}_N=\frac{\vec{p}_s\times\vec{p}_-}{|\vec{p}_s\times\vec{p}_-|}, \tag{23}$$

$$\hat{e}_T=\hat{e}_N\times\hat{e}_L, \tag{24}$$

with $\vec{p}_-$ and $\vec{p}_s$ being the three-momentum of l^- and s quark in the c.m. frame of l^+l^-.

The differential decay rate of $B\to X_s l^+l^-$ for any given spin direction $\hat{n}$ of lepton l^- may then be written as

$$\frac{dB(n)}{d\hat{s}}=\frac{1}{2}\left(\frac{dB}{d\hat{s}}\right)_{\text{unpol}}[1+(P_L\hat{e}_L+P_T\hat{e}_T+P_N\hat{e}_N)\hat{n}], \tag{25}$$

where P_L, P_T, and P_N are functions of $\hat{s}$ which gives longitudinal, transverse, and normal polarization components of polarization, respectively. The polarization component $P_i(i=L,T,N)$ is obtained by evaluating

$$P_i(\hat{s})=\frac{dB(\hat{n}=\hat{e}_i)/d\hat{s}-dB(\hat{n}=-\hat{e}_i)/d\hat{s}}{dB(\hat{n}=\hat{e}_i)/d\hat{s}+dB(\hat{n}=-\hat{e}_i)/d\hat{s}}. \tag{26}$$

The results obtained using the effective Hamiltonian (6) is

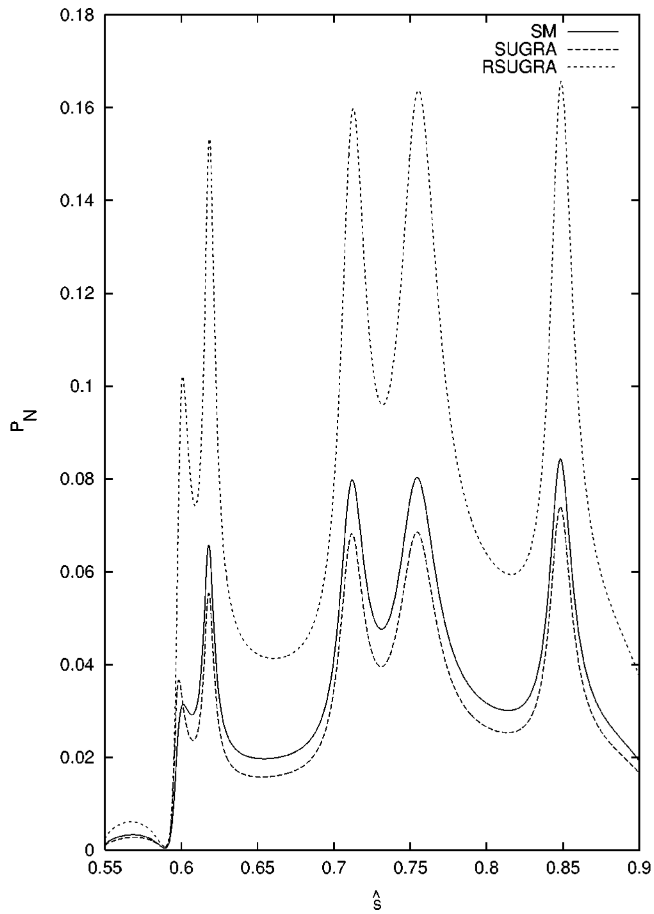

FIG. 5. Normal polarization asymmetry with $\hat{s}$. Other parameters are $\tan\beta = 30; m_0 = M = 130; A = -1$, for the relaxed SUGRA (RSUGRA) model $m_A = 120$. All masses are in GeV.

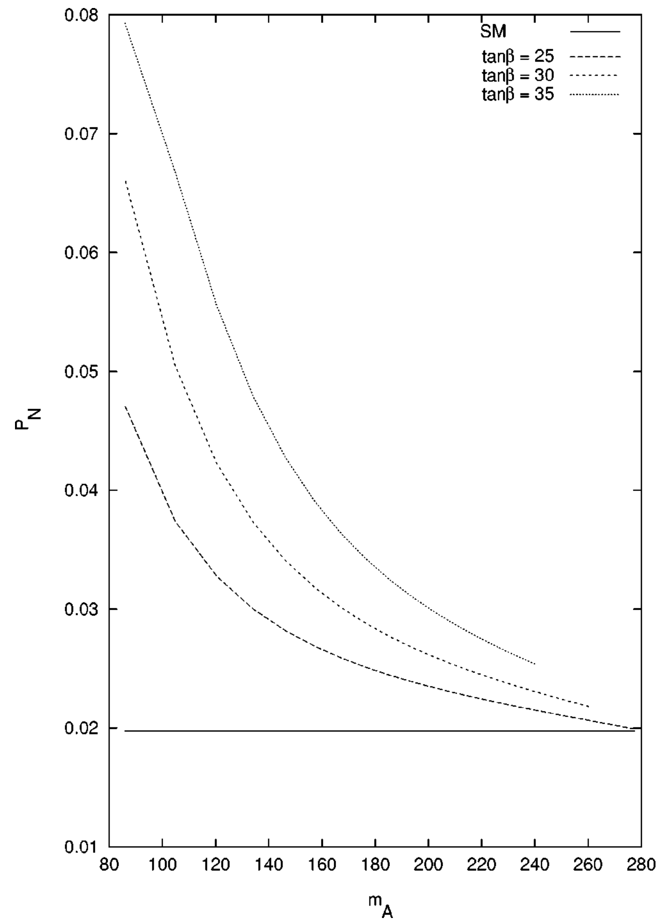

FIG. 6. Normal polarization asymmetry with m_A in relaxed SUGRA model. Other parameters are $m_0 = M = 130; A = -1; \hat{s} = 0.65$. All masses are in GeV.

$$P_L(\hat{s}) = \sqrt{1 - \frac{4\hat{m}_l^2}{\hat{s}}}[12C_7^{\rm eff}C_{10}[(1-\hat{m}_s^2)^2 - \hat{s}(1+\hat{m}_s^2)] + 2{\rm Re}(C_9^{\rm eff}C_{10})[(1-\hat{m}_s^2)^2 + \hat{s}(1+\hat{m}_s^2) - 2\hat{s}^2] + 6C_{Q_1}C_{10}\hat{m}_l(-1+\hat{m}_s^2+\hat{s}) + 3C_{Q_1}C_{Q_2} \times(-1-\hat{m}_s^2+\hat{s})\hat{s}]/\triangle, \tag{27}$$

$$P_T(\hat{s}) = \frac{3\pi\hat{m}_l}{2\sqrt{\hat{s}}}\lambda^{1/2}(1,\hat{s},\hat{m}_s^2)\Bigg[2C_7^{\rm eff}C_{10}(1-\hat{m}_s^2) - 4{\rm Re}(C_7^{\rm eff}C_9^{\rm eff})(1+\hat{m}_s^2) - \frac{4}{\hat{s}}|C_7^{\rm eff}|^2(1-\hat{m}_s^2)^2 + {\rm Re}(C_9^{\rm eff}C_{10})(1-\hat{m}_s^2) - |C_9^{\rm eff}|^2\hat{s} - \frac{1}{2}C_{Q_1}C_{10}\frac{4\hat{m}_l^2-\hat{s}}{\hat{m}_l} + C_{Q_2}C_7^{\rm eff}\frac{\hat{s}}{\hat{m}_l} + \frac{1}{2}{\rm Re}(C_{Q_2}C_9^{\rm eff})\frac{\hat{s}}{\hat{m}_l}\Bigg] \Bigg/ \triangle, \tag{28}$$

$$P_N(\hat{s}) = \frac{3\pi\hat{m}_l}{2\triangle}\sqrt{\hat{s}}\lambda^{1/2}(1,\hat{s},\hat{m}_s^2)\sqrt{1-\frac{4\hat{m}_l^2}{\hat{s}}}Im(C_9^{\rm eff*}) \times\left(\frac{1}{2}C_{Q_1} + C_{10}\hat{m}_l\right). \tag{29}$$

Expressions of P_L and P_T matches with Ref. [5] if C_{Q_1} and C_{Q_2} are absent, i.e., no NHB exchange effects. But P_T disagrees with Ref. [5] for a factor of 2 multiplying in term $C_7^* C_{10}$. Let us now focus our attention on the parameter space. Apart from gauge and Yukawa couplings, we have in the ''relaxed'' SUGRA model discussed above, six parameters m_0, M, Δ_0, A, B, and μ at the Planck scale. Use of renormalization group equations (RGE) allows one to evolve these parameters down to the electroweak scale M_W. At that scale the $SU(2)_L \times U(1)_Y$ spontaneously breaks down to $U(1)_{em}$ [Eq. (5)] ; v_1, v_2 are determined in the tree approximation by the Higgs boson potential with all its parameters scaled down to M_W. M_Z is related to v_1 and v_2 by

$$M_Z^2 = \frac{1}{2}(g^2 + g'^2)(v_1^2 + v_2^2), \tag{30}$$

with g,g' being, respectively, the $SU(2)_L$ and $U(1)_Y$ gauge couplings. Thus, for a given value of $\tan\beta = v_2/v_1$, and with all SM parameters given we have effectively four free parameters, which will be further subject to constraints arising out of the known limits on $b \to s\gamma$.

Figures 1–6 summarize our results, wherein we have presented the three polarization values in the SM, minimal SUGRA and the ''relaxed'' SUGRA (RSUGRA) as discussed before. The extra parameter in RSUGRA has been taken to be the CP-odd Higgs boson mass m_A which is related to the parameters in the potential by

$$m_A^2 = 2\Delta_0^2 + 2\mu^2 \tag{31}$$

with the parameters being evaluated at M_W. The general comment about all the three polarizations is that in SUGRA, there is no appreciable change from the SM value even with NHB contributions. This is because at high $\tan\beta$, the constraints obtained through $b \to s\gamma$ limits, forces the three neutral Higgs boson to large mass value thus suppressing the NHB contributions. This is precisely the reason that in relaxed version of SUGRA, where low Higgs boson mass become allowed, considerable deviations from SM values are possible.

Turning now to the absolute values of P_L, P_T, and P_N as shown in Figs. 1–6, it is important to note that at and around the resonant peaks, the dominant contributions come from the resonant B-W contributions, Eq. (7) multiplied by a phenomenologically empirical factor $\kappa_v = 2.35$. We have taken this factor to be universal for all resonances whereas the actual number is fitted only to J/ψ production. This introduces some uncertainty in values of the cross section around the higher resonances and it is for this reason that the polarization values given Figs. 1 and 6 are more reliable in between the $c\bar{c}$ resonances rather than at the resonances. Typically for $\tan\beta = 30$ in the region $0.63 \leqslant \hat{s} \leqslant 0.68$, as well as $0.77 \leqslant \hat{s} \leqslant 0.82$ the longitudinal polarization increases in magnitude by about 50%. A similar pattern occurs for P_T in the same region and in regions between higher resonances. For the normal polarization P_N in the two regions $0.63 \leqslant \hat{s} \leqslant 0.68$ and $0.77 \leqslant \hat{s} \leqslant 0.82$ the value changes by a factor of 2. In general in the regions between the resonances there are changes in the values of polarizations which are sufficiently large for experimental detection as and when data become available. Figures 2, 4, and 6 show the general dependence of the polarization parameters on $\tan\beta$ and m_A.

In conclusion our calculations indicate that in MSSM with a large $\tan\beta$ and low m_A value, the polarization asymmetries in $B \to X_s\tau^+\tau^-$ are sensitive to neutral Higgs boson exchange contributions. A similar kind of enhancements were also claimed in Ref. [14] but there the R-parity violating couplings were responsible for it, but here we are working in model where R parity is an exact symmetry. The usefulness of polarization measurements in the context of the standard model and beyond have already been emphasized in the literature [4,5,14] and our results are expected to be useful in comparing SUSY-model predictions with experimental results when they become available.

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