Dominance of Pion Exchange in R-Parity-Violating Supersymmetric Contributions to Neutrinoless Double Beta Decay

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We present a new contribution of the R-parity-violating ($\not{R}_p$) supersymmetry (SUSY) to neutrinoless double beta decay ($0\nu\beta\beta$) via the pion exchange between decaying neutrons. The pion coupling to the final state electrons is induced by the $\not{R}_p$ SUSY interactions. We have found this pion-exchange mechanism to dominate over the conventional two-nucleon one. The latter corresponds to direct interaction between quarks from two decaying neutrons without any light hadronic mediator like π meson. The constraints on the certain $\not{R}_p$ SUSY parameters are extracted from the current experimental $0\nu\beta\beta$-decay half-life limit. These constraints are significantly stronger than those previously known or expected from the ongoing accelerator experiments. [S0031-9007(96)02005-4]

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Neutrinoless double beta decay ($0\nu\beta\beta$) has long been recognized as a sensitive probe of the new physics beyond the standard model (SM) (see [1,2]). Various mechanisms of $0\nu\beta\beta$ decay were proposed and studied in the last two decades. The conventional mechanism is based on the exchange of a massive Majorana neutrino between the two decaying neutrons. A new mechanism was found within supersymmetric (SUSY) models with R-parity violation ($\not{R}_p$) in [3]. [$R_p = (-1)^{3B+L+2S}$ where S, B, and L are spin, baryon, and lepton numbers.] It was later studied in more details in [4]. A complete analysis of this mechanism within the minimal supersymmetric standard model (MSSM) was carried out in [5].

The nuclear $0\nu\beta\beta$ decay is triggered by the $0\nu\beta\beta$ quark transition $d + d \to u + u + 2e^-$, which is induced by certain fundamental interactions. It was a common practice to put the initial d quarks separately inside the two initial neutrons of a $0\nu\beta\beta$-decaying nucleus. This is the so called two-nucleon mode of the $0\nu\beta\beta$ decay [see Fig. 1(a)]. If the above $0\nu\beta\beta$-quark transition proceeds at short distances, as in the case of $\not{R}_p$ SUSY interactions, then the basic nucleon transition amplitude $n + n \to p + p + 2e^-$ is strongly suppressed for relative distances larger than the mean nucleon radius.

In this Letter we propose a new pion-exchange SUSY mechanism which is based on the double-pion exchange between the decaying neutrons [Fig. 1(b)]. At the quark level this mechanism implies the same short-distance $\not{R}_p$ MSSM interactions as in [5]. However, it essentially differs from the previous consideration of the SUSY contribution to the $0\nu\beta\beta$ decay at the stage of the hadronization. We assume that the $\not{R}_p$ MSSM quark interactions induce $\pi\pi \to 2e$ transition at the middle point of the diagram in Fig. 1(b). The importance of the pion-exchange currents in $0\nu\beta\beta$ decay was first pointed out by Pontecorvo [6]. Later, this idea was quantitatively realized in [7,8] for the case of the heavy Majorana neutrino exchange. It was shown that the pion-exchange contribution cannot be neglected in this case. We will show that in the case of the $\not{R}_p$ MSSM induced quark transition the pion-exchange contribution absolutely dominates over the conventional two-nucleon mode.

The R_p-violating part of the superpotential breaking lepton number conservation is

$$W_{\not{R}_p} = \lambda_{ijk} L_i L_j \overline{E}_k + \lambda'_{ijk} L_i Q_j \overline{D}_k . \qquad (1)$$

Here L, Q are lepton and quark doublets while $\overline{E}$, $\overline{D}$ are lepton and *down* quark singlet superfields. Indices i, j, k denote generations and $\lambda_{ijk} = -\lambda_{jik}$. In what follows we concentrate on the so called "direct" SUSY contribution to the $0\nu\beta\beta$ [3–5] depending only on the λ' term of the superpotential in Eq. (1). The combination of both λ' and λ terms may lead to the "indirect" SUSY contribution accompanied by the neutrino exchange [9], which we do not consider in the present Letter.

Starting from the λ' term in Eq. (1) the following effective quark-electron vertex has been derived [5]:

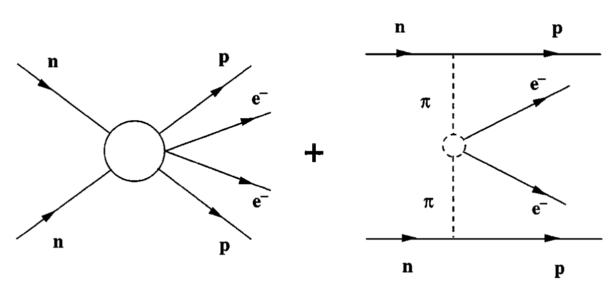

FIG. 1. (a) Two-nucleon mode $\mathcal{R}^{2N}_{0\nu\beta\beta}$ and (b) π-exchange $\mathcal{R}^{\pi N}_{0\nu\beta\beta}$ contributions to $0\nu\beta\beta$-decay matrix element $\mathcal{R}_{0\nu\beta\beta} = \mathcal{R}^{2N}_{0\nu\beta\beta} + \mathcal{R}^{\pi N}_{0\nu\beta\beta}$.

$$\mathcal{L}_{qe} = \frac{G_F^2}{2m_p}\,\overline{e}(1 + \gamma_5)e^c\left[(\eta_{\tilde{q}} + \eta_{\tilde{f}})(J_P J_P + J_S J_S) - \frac{1}{4}\,\eta_{\tilde{q}} J_T^{\mu\nu} J_{T\mu\nu})\right]. \quad (2)$$

These interactions violate the electron number $\Delta L_e = 2$. They are induced by heavy SUSY particles in a virtual intermediate state. An example of the Feynman diagram contributing to $\mathcal{L}_{qe}$ is given in Fig. 2. The color-singlet hadronic currents in Eq. (2) are $J_P = \overline{u}^\alpha \gamma_5 d_\alpha$, $J_S = \overline{u}^\alpha d_\alpha$, $J_T^{\mu\nu} = \overline{u}^\alpha \sigma^{\mu\nu}(1 + \gamma_5)d_\alpha$, where α is a color index.

The lepton number violating parameters η in Eq. (2) can be written in the following form

$$\eta_{\tilde{q}} = \Lambda^2\left[2\alpha_s\,\frac{m_p}{m_{\tilde{g}}} + \frac{3}{4}\,\alpha_2\,\frac{m_p}{m_\chi}\,(\epsilon_{Rd}^2 + \epsilon_{Lu}^2)\right], \quad (3)$$

$$\eta_{\tilde{f}} = \Lambda^2\left[2\alpha_s\,\frac{m_p}{m_{\tilde{g}}} + \frac{3}{2}\,\alpha_2\,\frac{m_p}{m_\chi}\left(\frac{m_{\tilde{q}}}{m_{\tilde{e}}}\right)^2 C\right]. \quad (4)$$

Here $\Lambda = (\sqrt{2\pi}/3)\lambda'_{111}G_F^{-1}m_{\tilde{q}}^{-2}$ and $C = 6(m_{\tilde{q}}/m_{\tilde{e}})^2 \times \epsilon_{Le}^2 - \epsilon_{Rd}\epsilon_{Le} - \epsilon_{Lu}\epsilon_{Rd}(m_{\tilde{e}}/m_{\tilde{q}})^2 - \epsilon_{Lu}\epsilon_{Le}$. $\alpha_2 = g_2^2/(4\pi)$ and $\alpha_s = g_3^2/(4\pi)$ are $\mathrm{SU}(3)_L$ and $\mathrm{SU}(3)_c$ gauge coupling constants; $m_{\tilde{g}}$ and m_χ are masses of the gluino $\tilde{g}$ and of the lightest neutralino χ. The latter is a linear combination of the gaugino and higgsino fields $\chi = \alpha_\chi \tilde{B} + \beta_\chi \tilde{W}^3 + \delta_\chi \tilde{H}_1^0 + \gamma_\chi \tilde{H}_2^0$. Here $\tilde{W}^3$ and $\tilde{B}$ are neutral $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$ gauginos while $\tilde{H}_2^0$, $\tilde{H}_1^0$ are higgsinos which are the superpartners of the two neutral Higgs boson fields H_1^0 and H_2^0 with a weak hypercharge $Y = -1, +1$, respectively. The mixing coefficients $\alpha_\chi, \beta_\chi, \gamma_\chi, \delta_\chi$ can be obtained from diagonalization of the 4×4 neutralino mass matrix. Neutralino couplings are defined as $\epsilon_{L\psi} = -T_3(\psi)\beta_\chi + \tan\theta_W[T_3(\psi) - Q(\psi)]\alpha_\chi$, $\epsilon_{R\psi} = Q(\psi)\tan\theta_W\alpha_\chi$ [10]. Here Q and T_3 are the electric charge and weak isospin of the fields $\psi = u, d, e$. In Eqs. (3) and (4) we used the universal squark mass $m_{\tilde{q}}$ ansatz at the weak scale $m_{\tilde{u}} \approx m_{\tilde{d}} \approx m_{\tilde{q}}$. This approximation is justified by the constraints from the flavor changing neutral currents and is sufficient for our analysis.

Now we have to reformulate the quark-lepton interactions in Eq. (2) in terms of the effective hadron-lepton interactions, which is necessary for the further nuclear structure calculations. The effective Lagrangian taking into account both the nucleon (p, n) and π-meson degrees of freedom in a nucleus can be written as follows:

$$\begin{aligned}\mathcal{L}_{he} &= \mathcal{L}_{2N} + \mathcal{L}_{\pi e} + \mathcal{L}_{\pi N} \\ &= \frac{G_F^2}{2m_p}\,\overline{p}\Gamma^{(i)}n\overline{p}\Gamma^{(i)}n\overline{e}(1 + \gamma_5)e^c \\ &\quad - \frac{G_F^2}{2m_p}\,m_\pi^4 a_\pi(\pi^-)^2\overline{e}(1 + \gamma_5)e^c \\ &\quad + g_s\overline{p}i\gamma_5 n\pi^+ .\end{aligned} \quad (5)$$

Here $\mathcal{L}_{2N}$, $\mathcal{L}_{\pi e}$, and $\mathcal{L}_{\pi N}$ describe the conventional two-nucleon mode, pion-exchange mode, and pion-nucleon interactions, respectively. They correspond to the first, second, and third terms of the second part of Eq. (5). $g_s = 13.4 \pm 1$ is known from experiment. The two-nucleon mode contributions $\mathcal{L}_{2N}$ to the $0\nu\beta\beta$ decay with different operator structures $\Gamma^{(i)}$ were derived and studied in [4,5] within the $\not{R}_p$ MSSM.

In this Letter we concentrate on the effect of the pion-exchange term $\mathcal{L}_{\pi e}$. The basic parameter a_π of the Lagrangian $\mathcal{L}_{\pi e}$ can be approximately related to the parameters of the fundamental Lagrangian $\mathcal{L}_{qe}$ using the on-mass-shell "matching condition" $\langle\pi^+|\mathcal{L}_{qe}|\pi^-\rangle = \langle\pi^+|\mathcal{L}_{\pi e}|\pi^-\rangle$. The solution of this equation is $a_\pi = \frac{1}{2}(\eta_{\tilde{q}} + \eta_{\tilde{f}})(c_P + c_S) - \frac{1}{8}\,\eta_{\tilde{q}}c_T$, where $\langle\pi^+|J_iJ_i| \times \pi^-\rangle = -m_\pi^4 c_i$ with $i = P, S, T$. Thus we obtain the approximate hadronic "image" $\mathcal{L}_{\pi e}$ of the fundamental quark-lepton Lagrangian $\mathcal{L}_{qe}$ given in Eq. (2).

The contribution of the $J_{P,S,T}$ currents to a_π can be estimated within the vacuum insertion approximation (VIA). Applying partial conservation of axial current (PCAC) we obtain

$$\begin{aligned}\langle\pi^+|J_PJ_P|\pi^-\rangle &= \frac{8}{3}\langle\pi^+|J_P|0\rangle\langle 0|J_P|\pi^-\rangle \\ &= -\frac{16}{3}\,f_\pi^2\,\frac{m_\pi^4}{(m_u + m_d)^2} \equiv -m_\pi^4 c_P\,,\end{aligned} \quad (6)$$

where $8/3$ is a combinatorial color factor and $f_\pi = 0.668\ m_\pi$. Taking the conventional values of the current quark masses $m_u = 4.2$ MeV, $m_d = 7.4$ MeV one gets $c_P \approx 342$. Within the VIA we have $c_s = c_T = 0$ since $\langle 0|J_S|\pi(p_\pi)\rangle = \langle 0|J_T^{\mu\nu}|\pi(p_\pi)\rangle = 0$. The scalar current matrix element vanishes due to the parity arguments, the tensor one vanishes due to $J_T^{\mu\nu} = -J_T^{\nu\mu}$ and the impossibility of constructing antisymmetric object having only one 4-vector p_π. Thus we expect the J_P contribution to be dominant.

The J_P dominance also follows from the nonrelativistic quark model (QM) [11]. Within this model one can calculate $\langle\pi^+|J_PJ_P|\pi^-\rangle$ using the closure approximation for the intermediate meson states [7]. After quite tedious calculations we end up again with $c_p \gg c_{S,T}$. In this case the numerical value $c_P \approx 1100$ is larger than that

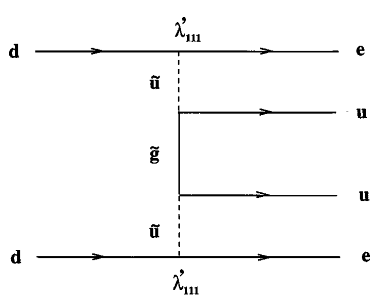

FIG. 2. An example of the supersymmetric contribution to $0\nu\beta\beta$ decay.

in Eq. (6), since in addition to the vacuum state there are other intermediate states taken into account. In what follows we use both the VIA and the QM values of c_P.

The large coefficient c_P enhances the pion-exchange contribution to the $0\nu\beta\beta$ decay. This enhancement factor is a generic property of the $R\!\!\!/_p$ SUSY models generating at low energies the $J_P J_P$ interactions [see Eq. (2)]. There is another factor enhancing the pion-exchange contribution compared to the two-nucleon mode. As explained later on, it stems from the fact that the pion exchange is longer ranged and thus covers a larger interval of the internucleon distances enhancing the nuclear matrix elements over the two-nucleon mode.

Starting from the Lagrangian $\mathcal{L}_{he}$ in Eq. (5) it is straightforward to calculate the contribution to the $0\nu\beta\beta$-matrix element $\mathcal{R}_{0\nu\beta\beta}$ which corresponds to the fundamental vertex $\mathcal{L}_{qe}$ in Eq. (2). It consists of the two terms $\mathcal{R}_{0\nu\beta\beta} = \mathcal{R}^{2N}_{0\nu\beta\beta} + \mathcal{R}^{\pi N}_{0\nu\beta\beta}$ describing the conventional two-nucleon mode $\mathcal{R}^{2N}_{0\nu\beta\beta}$ and the pion-exchange contribution $\mathcal{R}^{\pi N}_{0\nu\beta\beta}$. The relevant Feynman diagrams are given in Fig. 1. The corresponding half-life formula reads

$$[T^{0\nu\beta\beta}_{1/2}(0^+ \to 0^+)]^{-1} = G_{01}\left(\frac{m_A}{m_p}\right)^4 \times |\eta_{\tilde{q}}\mathcal{M}^{2N}_{\tilde{q}} + \eta_{\tilde{f}}\mathcal{M}^{2N}_{\tilde{f}} + (\eta_{\tilde{q}} + \eta_{\tilde{f}})\mathcal{M}^{\pi N}|^2 . \quad (7)$$

Here G_{01} is the standard phase space factor tabulated for various nuclei in [2] and $m_A = 850$ MeV. The analytic form of the two-nucleon mode nuclear matrix elements $\mathcal{M}^{2N}_{\tilde{q},\tilde{f}}$ are given in [5]. Here we present the new pion-exchange nuclear matrix element defined as

$$\mathcal{M}^{\pi N} = \frac{m_p}{m_e}\alpha^{\pi}(\mathcal{M}_{GT,\pi} + \mathcal{M}_{T,\pi}) . \quad (8)$$

The partial Gamow-Teller and tensor matrix elements are

$$\mathcal{M}_{GT,\pi} = \left\langle 0^+_f \left| \sum_{i\neq j}\tau^+_i\tau^+_j\sigma_{ij}\left(\frac{R}{r_{ij}}\right)F_1(x_\pi)\right| 0^+_i\right\rangle, \quad (9)$$

$$\mathcal{M}_{T,\pi} = \left\langle 0^+_f \left| \sum_{i\neq j}\tau^+_i\tau^+_j S_{ij}\left(\frac{R}{r_{ij}}\right)F_2(x_\pi)\right| 0^+_i\right\rangle, \quad (10)$$

where

$$S_{ij} = 3\vec{\sigma}_i\cdot\hat{\vec{r}}_{ij}\vec{\sigma}_j\cdot\hat{\vec{r}}_{ij} - \vec{\sigma}_i\cdot\vec{\sigma}_j, \quad \sigma_{ij} = \vec{\sigma}_i\cdot\vec{\sigma}_j ,$$
$$\hat{\vec{r}}_{ij} = (\vec{r}_i - \vec{r}_j)/|\vec{r}_i - \vec{r}_j|, \quad r_{ij} = |\vec{r}_i - \vec{r}_j| , \quad (11)$$

and $x_\pi = m_\pi r_{ij}$. Here $\vec{r}_i$ is the coordinate of the "ith" nucleon. The pion-nucleon structure coefficient in Eq. (8) is given by

$$\alpha^{\pi} = \frac{1}{96}\left(\frac{m_A}{m_p}\right)^2\left(\frac{m_\pi}{m_A}\right)^4\left(\frac{g_s}{f_A}\right)^2 c_P , \quad (12)$$

where $f_A = 1.261$. The pion-exchange SUSY potentials are

$$F_1(x) = (x - 2)e^{-x}, \qquad F_2(x) = (x + 1)e^{-x} . \quad (13)$$

The most stringent experimental lower limit on the $0\nu\beta\beta$-decay half-life has been obtained for ^{76}Ge [13], that favors especially this nucleus for nuclear structure calculations. In this Letter we direct our attention only to this isotope. We have employed the renormalized quasiparticle random phase approximation with proton-neutron pairing (full-RQRPA) [12] to calculate both the two-nucleon and pion-exchange nuclear matrix elements governing the $R\!\!\!/_p$ SUSY $0\nu\beta\beta$-decay of ^{76}Ge. The full-RQRPA includes the Pauli effect of fermion pairs and does not collapse for a physical value of the nuclear force strength. To include the Pauli principle more correctly we do not use the quasiboson approximation to derive the quasiparticle random phase approximation (QRPA). If one includes the exact Fermion commutation relations for nucleon pairs (two quasiparticles) as a QRPA expectation value, one obtains the renormalized QRPA (RQRPA), which is stable against the collapse of $2\nu\beta\beta$ Gamow-Teller transition. Therefore the RQRPA offers a significantly more reliable treatment of the nuclear many-body problem for the description of the $0\nu\beta\beta$ decay. Thus it also allows one to establish more reliable constraints on the $R\!\!\!/_p$ SUSY parameters from the best available experimental lower bound on the $0\nu\beta\beta$-decay half-life. We have found the following numerical values of the nuclear matrix elements for ^{76}Ge: $\mathcal{M}^{2N}_{\tilde{q}} = -61$; $\mathcal{M}^{2N}_{\tilde{f}} = 0.85$; $\mathcal{M}^{\pi N} = -1800$(QM), -600(VIA). The pion-exchange matrix element is given for QM and VIA values of the coefficient c_P. It is apparent that in both QM and VIA cases the dominant contribution to Eq. (7) comes from the pion-exchange mechanism corresponding to $\mathcal{M}^{\pi N}$. The VIA value $\mathcal{M}^{\pi N} = -600$ we will use for conservative estimations. It is worthwhile to note that the above nuclear matrix elements are quite stable with respect to variation of the nuclear model parameters. The uncertainty of the calculated values of $\mathcal{M}^{2N}_{\tilde{q},\tilde{f}}$ and $\mathcal{M}^{\pi N}$ does not exceed 20%.

Now we are ready to extract the constraints on the $R\!\!\!/_p$ MSSM parameters from the nonobservation of $0\nu\beta\beta$ decay. The current experimental lower bound on the ^{76}Ge $0\nu\beta\beta$-decay half-life [13] is $T^{0\nu\beta\beta-\mathrm{exp}}_{1/2}(0^+ \to 0^+) \geq 9.1 \times 10^{24}$ years 90% c.l. Combining this bound with Eq. (7) and the above given numerical values of $\mathcal{M}^{\pi N}$ we get a constraint on the sum of the effective MSSM parameters $\eta_{\tilde{q}} + \eta_{\tilde{f}} \leq 2.1 \times 10^{-9}$. If one does not include the pion-exchange contribution then one gets a constraint $\eta_{\tilde{q}} + 0.014\eta_{\tilde{f}} \leq 7.8 \times 10^{-8}$ from the remaining $2N$ mode. It is essentially less stringent than the above given πN mode constraint by more than 1 order of magnitude for $\eta_{\tilde{q}}$ and by 3 orders for $\eta_{\tilde{f}}$.

The gluino and neutralino contributions to η_i cannot cancel each other within the present experimental limits on their masses and couplings (see [5]). Therefore we can extract from the above limit on η_i the constraints on these individual contributions. The gluino contribution constraint is

$$\lambda'_{111} \leq 2.0(1.18) \times 10^{-4}\left(\frac{m_{\tilde{q}}}{100\ \text{GeV}}\right)^2\left(\frac{m_{\tilde{g}}}{100\ \text{GeV}}\right)^{1/2} \tag{14}$$

for the VIA(QM) value of the pion matrix element parameter c_P. The neutralino contribution constraint is more complex because it involves more parameters; neutralino mixing coefficients, selectron, and squark masses. However, it can be cast into the form of Eq. (14) under the phenomenologically viable simplifying assumptions. Assume that the neutralino is B-ino dominant $\alpha_\chi \gg \beta_\chi, \delta_\chi, \gamma_\chi$ and that $m_{\tilde{q}} \geq m_{\tilde{e}}$. Then we get

$$\lambda'_{111} \leq 5.2(3.07) \times 10^{-4}\left(\frac{m_{\tilde{e}}}{100\ \text{GeV}}\right)^2\left(\frac{m_{\chi}}{100\ \text{GeV}}\right)^{1/2}. \tag{15}$$

If all SUSY particle masses in Eqs. (14) and (15) were at their present experimental lower bounds [14] $m_{\tilde{q}} \geq 90$ GeV, $m_{\tilde{e}} \geq 45$ GeV, $m_{\tilde{g}} \geq 100$ GeV, $m_\chi \geq 19$ GeV, we could estimate the size of the $R\!\!\!/_p$ coupling constant $\lambda'_{111} \leq 4.6(2.7) \times 10^{-5}$. A conservative bound can be obtained using the SUSY 'naturalness" upper bound $m_{\tilde{g},\tilde{q},\chi} \leq 1$ TeV. It gives $\lambda'_{111} \leq 6.3(3.7) \times 10^{-2}$. This limit and those in Eqs. (14) and (15) are the best known limits on the $R\!\!\!/_p$ coupling λ'_{111} [see [5] and references therein].

It is interesting to compare the $0\nu\beta\beta$ decay and accelerator experiments from the point of view of their sensitivity to the $R\!\!\!/_p$ SUSY signal. Previously [5] it was shown that even in the $2N$ mode the constraints from the $0\nu\beta\beta$ decay exclude the domain of the $R\!\!\!/_p$ MSSM parameter space accessible for the ongoing experiments with the ZEUS detector at HERA. This conclusion touched upon the region of the so called $R\!\!\!/_p$ resonant single squark $\tilde{q}$ production mechanism in deep inelastic ep-scattering [15]. Taking into account the πN mode makes this conclusion much stronger. The reason is that the excluded region becomes significantly larger than in case of the $2N$ mode alone. Now this region extends so far that it would include the HERA domain even if the experimental lower bound on $T_{1/2}^{0\nu\beta\beta}$ was by about (conservatively) 5 orders of magnitude less than the currently existing one.

Nevertheless, there is still a window for the HERA experiments in the region corresponding to another mechanism assuming R_p-conserving $\tilde{e} + \tilde{q}$ production and their subsequent $R\!\!\!/_p$ cascade decays. This region is $m_{\tilde{e}} + m_{\tilde{q}} \leq 205$ GeV and $\lambda'_{111} \geq 10^{-6}$ [16]. [Note that unlike $0\nu\beta\beta$-decay searches HERA can probe within this mechanism $R\!\!\!/_p$ coupling constants other than λ'_{111}.] In the present Letter we would like to stress that the $0\nu\beta\beta$-decay constraints given in Eqs. (14) and (15) dramatically reduces the above quoted regions. Only a narrow part of it corresponding to a very low values of the $R\!\!\!/_p$ coupling constant $10^{-6} \leq \lambda'_{111} \leq 0.76(2.4) \times 10^{-4}(\frac{m_{\tilde{e}}+m_{\tilde{q}}}{100\ \text{GeV}})^2$ and $m_{\tilde{e}} + m_{\tilde{q}} \leq 205$ GeV remains not excluded by the $0\nu\beta\beta$-decay constraints. Here we used the conservative VIA value of the pion matrix element c_P and put $m_\chi = m_{\tilde{g}} = 100$ GeV (1 TeV).

Summarizing, we point out that the SUSY contribution to the $0\nu\beta\beta$ decay comes dominantly via and pion-exchange mechanism considered in the present letter. The conventional two-nucleon mechanism [3–5], corresponding to $nn \rightarrow ppee$ transition without light particles (pion or neutrino) in the intermediate state, brings only a subdominant SUSY contribution. In practice the pion-exchange mechanism considerably enhances the sensitivity of the $0\nu\beta\beta$ decay to the supersymmetry. This allowed us to obtain presently the most stringent limitations of the certain first generation $R\!\!\!/_p$ MSSM parameters.

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