

**Direct detection of the Wino and Higgsino-like neutralino dark matter at one-loop level**Junji Hisano,<sup>1</sup> Shigeki Matsumoto,<sup>1</sup> Mihoko M. Nojiri,<sup>2</sup> and Osamu Saito<sup>1</sup><sup>1</sup>*ICRR, University of Tokyo, Kashiwa 277-8582, Japan*<sup>2</sup>*YITP, Kyoto University, Kyoto 606-8502, Japan*

(Received 23 July 2004; published 19 January 2005)

The neutralino-nucleon ( $\tilde{\chi}^0-N$ ) scattering is an important process for direct dark matter searches. In this paper we discuss one-loop contributions to the cross section in the winolike and Higgsino-like LSP cases. The neutralino-nucleon scattering mediated by the Higgs  $\tilde{\chi}^0\tilde{\chi}^0$  and  $Z\tilde{\chi}^0\tilde{\chi}^0$  couplings at tree level is suppressed by the gaugino-Higgsino mixing at tree-level when the neutralino is close to a weak eigenstate. The one-loop contribution to the cross section, generated by the gauge interaction, is not suppressed by any SUSY particle mass or mixing in the wino- and Higgsino-like LSP cases. It may significantly alter the total cross section when  $\sigma_{\tilde{\chi}^0 N} \sim 10^{-45}$  cm<sup>2</sup> or less.

DOI: 10.1103/PhysRevD.71.015007

PACS numbers: 14.80.Ly, 95.35.+d

**I. INTRODUCTION**

Dark matter mass density in the Universe is now measured very precisely by cosmological observations,  $\Omega_M = 0.27 \pm 0.04$  [1,2]. Now one of the important questions regarding to the dark matter is the constituent. The minimal supersymmetric standard model (MSSM) predicts the stable lightest supersymmetric particle (LSP) if the R parity is conserved. This is one of the attractive features of the model, because the neutralino LSP is a good candidate of the dark matter in the Universe because it is a weakly-interacting massive particle (WIMP).

The neutralino LSP is a linear combination of gauginos (bino and wino) and Higgsinos, which are superpartners of the gauge and Higgs bosons in the SM, respectively. The binolike or Higgsino-like neutralino LSP is the dark matter candidate in the minimal supergravity model (MSUGRA). Recently many authors have investigated the cosmological relic density of the binolike dark matter in the MSUGRA. It is found that the thermal relic density of the LSP is too large compared to the current observations, unless the coannihilations with other SUSY particles enhance the effective neutralino annihilation cross section at the early universe or entropy production reduces the number density after the decoupling of the LSP. The anomaly mediated SUSY-breaking model [3,4] and the string models with moduli dominated SUSY-breaking [5] predict the winolike or Higgsino-like neutralino LSP. The thermal relic density of the neutralino LSP is too low in the models unless the LSP mass ( $m_{\tilde{\chi}^0}$ ) is above 1 TeV. However, decay of gravitino or other quasistable particles may produce the dark matter nonthermally so that the relic abundance is consistent with the observation [6,7]. Also, while the LSP with the mass heavier than about 1 TeV may lead to the naturalness problem, it may be consistent in the split SUSY model [8].

Many experiments are now searching for the direct or indirect evidence of the dark matter. The counting rates in the direct search experiments depend on the LSP neutralino-nucleon ( $\tilde{\chi}^0-N$ ) interactions. The cross section above  $10^{-42}$  cm<sup>2</sup> is now explored for  $m_{\tilde{\chi}^0} \lesssim 1$  TeV. While

the annual modulation observed by the DAMA experiment corresponds to the  $\tilde{\chi}^0-N$  spin-independent cross section around  $10^{-42}$  cm<sup>2</sup> [9], the recent result of the CDMSII rejects whole of the DAMA signal region [10] if the spin-dependent part of the interaction is negligible. The sensitivity to the dark matter signal may be improved up to  $10^{-(45-46)}$  cm<sup>2</sup> or more in future.

The  $\tilde{\chi}^0-N$  scattering cross section is sensitive to nature of the LSP and the SUSY particle mass spectrum. The Higgs and Z boson exchanges are dominant contributions to the spin-independent and spin dependent interactions responsible to the scattering, respectively, in the wide parameter region. They are suppressed by the gaugino-Higgsino mixing at tree level. When gauginos or Higgsino is much heavier than the weak scale, the LSP is close to a pure weak eigenstate and the scattering cross section is strongly suppressed. The squark exchange also contributes to the cross section, and it is not suppressed by the mixing. However, it tends to be subdominant due to the heavier squark masses in the typical models.

In this paper, we evaluate the one-loop radiative corrections to the wino- and Higgsino-like LSP scattering cross sections on a nucleon, which are induced by the gauge interaction. These one-loop corrections to the cross section are not suppressed by the Higgsino-gaugino mixing. In addition to it, the loop integrals are only suppressed by the weak gauge boson masses, because the chargino, which is the  $SU(2)$  partner of the LSP, is degenerated with the LSP in mass. The gauge-loop correction can dominate the total cross section in a limit that the LSP is almost a pure weak eigenstate, setting the “lower limit” of the total scattering cross section. When the LSP is binolike, the one-loop correction is negligible since bino does not have gauge charges.

We note that similar phenomena sometimes appear in radiative correction to the SUSY processes. For example, it is known that the mass difference between the LSP and chargino may be dominated by the radiative correction due to the gauge loops in the case of the winolike LSP. The tree-level mass difference is  $O(m_Z^4/M_{SUSY}^3)$ . On the other

hand, the radiative one is not suppressed by any SUSY particle mass, and it is proportional to  $\alpha_2 m_W$ . This is because the loop momentum around  $m_W$  dominates in the loop integrals due to the mass degeneracy between the LSP and the chargino. In addition, the pair annihilation cross sections of the wino- and Higgsino-like neutralino LSPs to two gammas at one-loop level are not suppressed by the neutralino mass [11]. It is rather enhanced by a nonperturbative effect when the mass is larger than the  $m_W/\alpha_2$  [12]. This effect also comes from the mass degeneracy between the LSP and the chargino.

This paper is organized as follows. In Sec. II, we set up the formula for the general low-energy effective action and present the  $\tilde{\chi}^0-N$  scattering cross section. The interactions are classified into the spin-dependent and spin-independent ones at the nonrelativistic limit of the neutralino LSP. We note that there are two classes for the spin-independent contributions; one is proportional to the scalar operator of quark,  $\langle N|m_q\bar{q}q|N\rangle$ , and the other is proportional to the twist-2 operator,  $\langle N|\bar{q}i(\partial_\mu\gamma_\nu + \partial_\nu\gamma_\mu - 1/2g_{\mu\nu}\not{\partial})q|N\rangle$ . The tree-level contribution to the twist-2 operators, which is induced by the squark exchange, is negligible [13]. However, the twist-2 operators is sizable at one-loop level.

In Sec. III, we briefly summarize the tree-level  $\tilde{\chi}^0$  couplings responsible to the  $\tilde{\chi}^0-N$  scattering for the wino- and Higgsino-like LSPs. In Sec. IV, the gauge-loop correction to the effective action in the winolike LSP case is presented. The dominant contribution comes from the  $W$  boson loops in this case. We identify the sources of the one-loop corrections to the spin-independent interaction, correction to the Higgs coupling of the LSP and those to scalar and twist-2 operators. The corrections are only suppressed by one-loop factors, not by the Higgsino-gaugino mixing nor any SUSY particle mass. The one-loop correction in the Higgsino-like LSP case is summarized in the Appendix since the structure of the radiative correction is similar to the winolike case.

In Sec. V, we present some numerical results. Among the gauge-loop contributions, the sign of the correction to the twist-2 operator is opposite to those to the scalar operator and Higgs boson vertex in the case of the winolike LSP. The cancellation reduces the total correction to the spin-independent cross section in wide region of the MSSM parameter space. Because of that, the spin-independent cross section induced by the gauge-loop alone is small, only around  $\sim 10^{-(46-47)}$  cm<sup>2</sup> in the limit when the tree-level contribution to the cross section is negligible. The total cross section, including the contributions at tree and one-loop levels, will be affected by the gauge-loop corrections when the cross section is close to the sensitivities of the proposed dark matter search experiments ( $\sigma_{\tilde{\chi}^0 N} \sim 10^{-(45-46)}$  cm<sup>2</sup>). For the Higgsino-like LSP, the cross section induced by the gauge-loop diagrams alone is 1 order of magnitude smaller than that of the winolike LSP. Section VI is devoted to conclusion and discussion.

## II. EFFECTIVE LAGRANGIAN FOR NEUTRALINO AND NUCLEON SCATTERING

In this section we present the effective Lagrangian for the  $\tilde{\chi}^0-N$  scattering and the cross section formula. The effective interactions of the neutralino LSP with light quarks and gluon at the renormalization scale  $\bar{\mu}_0 \simeq m_p$  are given in a limit of the nonrelativistic neutralino as follows

$$\mathcal{L}^{\text{eff}} = \sum_{q=u,d,s} \mathcal{L}_q^{\text{eff}} + \mathcal{L}_g^{\text{eff}}, \quad (1)$$

where

$$\begin{aligned} \mathcal{L}_q^{\text{eff}} = & d_q \bar{\chi}^0 \gamma^\mu \gamma_5 \tilde{\chi}^0 \bar{q} \gamma_\mu \gamma_5 q + f_q m_q \bar{\chi}^0 \tilde{\chi}^0 \bar{q} q \\ & + \frac{g_q^{(1)}}{m_{\tilde{\chi}^0}} \bar{\chi}^0 i \partial^\mu \gamma^\nu \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^q \\ & + \frac{g_q^{(2)}}{m_{\tilde{\chi}^0}^2} \bar{\chi}^0 (i \partial^\mu)(i \partial^\nu) \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^q, \end{aligned} \quad (2)$$

$$\begin{aligned} \mathcal{L}_g^{\text{eff}} = & f_G \bar{\chi}^0 \tilde{\chi}^0 G_{\mu\nu}^a G^{a\mu\nu} + \frac{g_G^{(1)}}{m_{\tilde{\chi}^0}} \bar{\chi}^0 i \partial^\mu \gamma^\nu \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^g \\ & + \frac{g_G^{(2)}}{m_{\tilde{\chi}^0}^2} \bar{\chi}^0 (i \partial^\mu)(i \partial^\nu) \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^g. \end{aligned} \quad (3)$$

Here, we include terms up to the second derivative of the neutralino field. The first term of  $\mathcal{L}_q^{\text{eff}}$  is a spin-dependent interaction, while the other terms in  $\mathcal{L}_q^{\text{eff}}$  and  $\mathcal{L}_g^{\text{eff}}$  are spin-independent ‘‘coherent’’ contributions. The third and fourth terms in  $\mathcal{L}_q^{\text{eff}}$  and the second and third terms in  $\mathcal{L}_g^{\text{eff}}$  depend on the twist-2 operators (traceless part of the energy momentum tensor) for quarks and gluon,

$$\begin{aligned} \mathcal{O}_{\mu\nu}^q & \equiv \frac{1}{2} \bar{q} i \left( \partial_\mu \gamma_\nu + \partial_\nu \gamma_\mu - \frac{1}{2} g_{\mu\nu} \not{\partial} \right) q, \\ \mathcal{O}_{\mu\nu}^g & \equiv \left( G_\mu^{a\rho} G_{\rho\nu}^a + \frac{1}{4} g_{\mu\nu} G_{\alpha\beta}^a G^{a\alpha\beta} \right). \end{aligned} \quad (4)$$

The scattering cross section of the neutralino with target nuclei is expressed compactly by using the coefficients given in  $\mathcal{L}_q^{\text{eff}}$  and  $\mathcal{L}_g^{\text{eff}}$  as follows [14],

$$\begin{aligned} \sigma = & \frac{4}{\pi} \left( \frac{m_{\tilde{\chi}^0} m_T}{m_{\tilde{\chi}^0} + m_T} \right)^2 \left[ (n_p f_p + n_n f_n)^2 \right. \\ & \left. + 4 \frac{J+1}{J} (a_p \langle S_p \rangle + a_n \langle S_n \rangle)^2 \right]. \end{aligned} \quad (5)$$

The first term in the bracket comes from the spin-independent interactions while the second one is generated by the spin-dependent one.

In Eq. (5),  $m_T$  is the target nucleus mass and  $n_p$  and  $n_n$  are proton and neutron numbers in the target nucleus, respectively. The spin-independent coupling of the neutralino with nucleon,  $f_N (N = p, n)$ , in Eq. (5) is given as

$$\begin{aligned}
f_N/m_N &= \sum_{q=u,d,s} \left[ f_q f_{Tq} + \frac{3}{4} [q(2, \bar{\mu}_0^2) + \bar{q}(2, \bar{\mu}_0^2)] \right. \\
&\quad \times (g_q^{(1)} + g_q^{(2)})(\bar{\mu}_0) \left. \right] - \frac{8\pi}{9\alpha_s} f_{TG} f_G \\
&\quad + \frac{3}{4} G(2, \bar{\mu}_0^2) (g_G^{(1)} + g_G^{(2)})(\bar{\mu}_0), \tag{6}
\end{aligned}$$

where the matrix elements of nucleon are expressed as

$$\begin{aligned}
f_{Tq} &\equiv \langle N | m_q \bar{q} q | N \rangle / m_N, \quad f_{TG} \\
&\equiv 1 - \sum_{u,d,s} f_{Tq}, \\
\langle N(p) | \mathcal{O}_{\mu\nu}^q | N(p) \rangle &= \frac{1}{m_N} \left( p_\mu p_\nu - \frac{1}{4} m_N^2 g_{\mu\nu} \right) \\
&\quad \times [q(2, \bar{\mu}_0^2) + \bar{q}(2, \bar{\mu}_0^2)], \tag{7} \\
\langle N(p) | \mathcal{O}_{\mu\nu}^g | N(p) \rangle &= \frac{1}{m_N} \left( p_\mu p_\nu - \frac{1}{4} m_N^2 g_{\mu\nu} \right) G(2, \bar{\mu}_0).
\end{aligned}$$

Here,  $q(2, \bar{\mu}_0^2)$ ,  $\bar{q}(2, \bar{\mu}_0^2)$  and  $G(2, \bar{\mu}_0^2)$  are the second moments of the quark, antiquark and gluon distribution functions, respectively,

$$\begin{aligned}
q(2, \bar{\mu}_0^2) + \bar{q}(2, \bar{\mu}_0^2) &= \int_0^1 dx x [q(x, \bar{\mu}_0^2) + \bar{q}(x, \bar{\mu}_0^2)], \\
G(2, \bar{\mu}_0^2) &= \int_0^1 dx x g(x, \bar{\mu}_0^2). \tag{8}
\end{aligned}$$

The constant  $a_N$  ( $N = p, n$ ), which is responsible for the spin-dependent contribution, is defined as

$$a_N = \sum_{q=u,d,s} d_q \Delta q_N, \tag{9}$$

$$2s_\mu \Delta q_N \equiv \langle N | \bar{q} \gamma_\mu \gamma_5 q | N \rangle, \tag{10}$$

where  $s_\mu$  is the nucleon's spin, while  $\langle S_N \rangle = \langle A | S_N | A \rangle$  in Eq. (5) is the expectation value of the third component of the spin operator of the proton or neutron group in the nucleus  $A$ .

The above formula does not contain heavy quark contributions explicitly. After integrating out the Higgs and weak gauge bosons and squarks, the effective interactions of the neutralino with the heavy quarks ( $Q$ ), which contribute to the spin-independent interaction, are

$$\begin{aligned}
\mathcal{L}_Q^{\text{eff}} &= f_Q m_Q \bar{\chi}^0 \tilde{\chi}^0 \bar{Q} Q + f'_Q \bar{\chi}^0 \tilde{\chi}^0 \bar{Q} i \not{\partial} Q \\
&\quad + \frac{g_Q^{(1)}}{m_{\tilde{\chi}^0}} \bar{\chi}^0 i \partial^\nu \gamma^\mu \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^Q \\
&\quad + \frac{g_Q^{(2)}}{m_{\tilde{\chi}^0}^2} \bar{\chi}^0 (i \partial^\mu) (i \partial^\nu) \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^Q. \tag{11}
\end{aligned}$$

The effective interactions (11) contribute to  $\mathcal{L}_g^{\text{eff}}$  through the heavy quark loop diagrams.

It is well-known that the matrix element for  $m_Q \bar{Q} Q$  in Eq. (11) can be replaced by that for  $-\alpha_s/(12\pi) G^{a\mu\nu} G_{\mu\nu}^a$  due to the trace anomaly. Then,

$$\langle N | m_Q \bar{Q} Q | N \rangle = \frac{2}{27} f_{TG} m_N. \tag{12}$$

The second term of Eq. (11) is proportional to  $\bar{Q} i \not{\partial} Q$ . It would reduce to  $m_Q \bar{Q} Q \bar{\chi}^0 \tilde{\chi}^0$ , if the equation of motion for the heavy quark could be applied. However, this is not justified because the heavy quark loop diagram induced by the interaction, which contributes to the operator  $\bar{\chi}^0 \tilde{\chi}^0 GG$ , has an UV divergence. Indeed it is found in Ref. [13], where the squark exchange contribution to the elastic scattering is evaluated, that an estimation of that operator using the equation of motion for the heavy quarks disagrees with the explicit full one-loop calculation by a factor of 2. We need to calculate the vertex for  $\bar{\chi}^0 \tilde{\chi}^0 GG$ , which is induced by the heavy quark loop diagrams, in the original theory. In this paper we parameterize the matrix element of  $\bar{Q} i \not{\partial} Q$  as

$$\langle N | \bar{Q} i \not{\partial} Q | N \rangle = a_{\text{eff}} \frac{2}{27} f_{TG} m_N. \tag{13}$$

by introducing a phenomenological parameter  $a_{\text{eff}}$ . The precise determination of  $a_{\text{eff}}$  requires evaluation of the higher-order loop diagrams, and it is out of scope of this paper.

Equations (2), (3), and (11) contain the traceless parts of the energy momentum tensor,  $\mathcal{O}_{\mu\nu}^q$ ,  $\mathcal{O}_{\mu\nu}^Q$  and  $\mathcal{O}_{\mu\nu}^g$ , whose matrix elements are scale-dependent. The second moments of the quark and antiquark distribution functions,  $q(2, \bar{\mu}^2)$  and  $\bar{q}(2, \bar{\mu}^2)$ , are mixed with that of the gluon distribution function,  $G(2, \bar{\mu}^2)$ , once the QCD radiative corrections are included. Their scale dependences are compensated by those of the coefficients  $g_q^{(1,2)}$  and  $g_G^{(1,2)}$ , so that the total cross section is scale-independent. Thus,

$$\begin{aligned}
&\sum_{q=u,d,s} [q(2, \bar{\mu}_0^2) + \bar{q}(2, \bar{\mu}_0^2)] (g_q^{(1)} + g_q^{(2)})(\bar{\mu}_0) \\
&\quad + G(2, \bar{\mu}_0^2) (g_G^{(1)} + g_G^{(2)})(\bar{\mu}_0) \\
&= \sum_{m_q \leq \bar{\mu}} [q(2, \bar{\mu}^2) + \bar{q}(2, \bar{\mu}^2)] (g_q^{(1)} + g_q^{(2)})(\bar{\mu}) \\
&\quad + G(2, \bar{\mu}^2) (g_G^{(1)} + g_G^{(2)})(\bar{\mu}). \tag{14}
\end{aligned}$$

In this paper we use the second moments for gluon and quark distribution functions at  $\bar{\mu} = m_Z$ , which are derived by the CTEQ parton distribution [15].

In Table I, we show the parameters for the matrix elements, used in this paper. The second moments for the up and down quark distribution functions are sizable. We will see in Sec. III the term proportional to  $q(2, \bar{\mu}^2)$  is generated by the one-loop box correction involving the  $W$  boson exchanges. The  $f_{Tq}$  and  $f_{TG}$  represent the fractions of the trace part of the energy momentum tensor, or the

quark and gluino contributions to the nucleon mass, as defined as above. The strange quark contribution to the spin-independent cross section is dominant among those for light quarks. Also, the term proportional to  $f_{TG}$  in Eq. (6) through heavy quark loops is not negligible.

### III. LSP ELASTIC SCATTERING INDUCED BY TREE-LEVEL LSP COUPLINGS

The mass matrix of the the neutralinos and charginos are given by

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -m_Z s_W \cos\beta & m_Z s_W \sin\beta \\ 0 & M_2 & m_Z c_W \cos\beta & -m_Z c_W \sin\beta \\ -m_Z s_W \cos\beta & m_Z c_W \cos\beta & 0 & -\mu \\ m_Z s_W \sin\beta & -m_Z c_W \sin\beta & -\mu & 0 \end{pmatrix} \quad (15)$$

$$\mathcal{M}_C = \begin{pmatrix} M_2 & \sqrt{2}m_W \sin\beta \\ \sqrt{2}m_W \cos\beta & \mu \end{pmatrix}, \quad (16)$$

which is written by the  $(\tilde{B}, \tilde{W}^0, \tilde{H}_1^0, \tilde{H}_2^0)$  bases and  $(\tilde{W}^+, \tilde{H}^+)$  bases, respectively.

When the lightest neutralino is winolike ( $M_2 \ll \mu, M_1$ ) or Higgsino-like ( $\mu \ll M_1, M_2$ ), the mass eigenvalues of the lightest neutralino and chargino are close to each other. They are given as

$$\begin{aligned} m_{\tilde{\chi}^0} &= M_2 + \frac{m_W^2}{M_2^2 - \mu^2}(M_2 + \mu \sin 2\beta) + \dots \\ m_{\tilde{\chi}^-} &= M_2 + \frac{m_W^2}{M_2^2 - \mu^2}(M_2 + \mu \sin 2\beta) + \dots \end{aligned} \quad (17)$$

TABLE I. Parameters for quark and gluon matrix elements used in this paper.  $f_{Ti}$  ( $i = u, d, s$ ) is taken from the estimation in Refs. [16–19]. The second moments for gluon and quarks at  $\bar{\mu} = m_Z$  and the spin fraction are for proton. Those for neutron are given by exchange of up and down quarks in the tables. The second moments are calculated using the CTEQ parton distribution [15].

For proton			
$f_{Tu}$			0.023
$f_{Td}$			0.034
$f_{Ts}$			0.14
For neutron			
$f_{Tu}$			0.019
$f_{Td}$			0.041
$f_{Ts}$			0.14
Spin fraction			
$\Delta u$			0.77
$\Delta d$			-0.49
$\Delta s$			-0.15
Second moment at $\bar{\mu} = m_Z$			
$G(2)$	0.48		
$u(2)$	0.22	$\bar{u}(2)$	0.034
$d(2)$	0.11	$\bar{d}(2)$	0.036
$s(2)$	0.026	$\bar{s}(2)$	0.026
$c(2)$	0.019	$\bar{c}(2)$	0.019
$b(2)$	0.012	$\bar{b}(2)$	0.012

for the wino case, and

$$\begin{aligned} m_{\tilde{\chi}^0} &= \mu + \frac{m_Z^2(1 + \sin 2\beta)}{2(\mu - M_1)(\mu - M_2)}(\mu - M_1 c_W^2 - M_2 s_W^2) \\ &+ \dots \end{aligned} \quad (18)$$

$$m_{\tilde{\chi}^-} = \mu - \frac{m_W^2}{M_2^2 - \mu^2}(\mu + M_2 \sin 2\beta) + \dots$$

for the Higgsino case ( $\mu > 0$ ). The lightest neutralino mass eigenstate becomes very close to the pure wino or Higgsino in the limit, therefore  $m_{\tilde{\chi}^0} \sim M_2$  (wino-like) or  $\sim \mu$  (Higgsino-like). The mass difference between the LSP and the lighter chargino  $\delta_c = (m_{\tilde{\chi}^-} - m_{\tilde{\chi}^0})/m_{\tilde{\chi}^0}$  is also very small if  $|\mu - M_2| \gg m_Z$ . The LSP and the lighter chargino form an  $SU(2)$  triplet state in the wino-like LSP case or vectorlike doublets in the Higgsino-like one with the second-lightest neutralino, respectively.

In this section, we discuss the tree-level contributions to the effective interactions in Eqs. (2), (3), and (11) in the wino-like and Higgsino-like LSP cases. The  $\tilde{\chi}^0-N$  spin-independent interactions are generated by the  $t$ -channel exchange of one Higgs boson and the  $s$ -channel exchange of one squark, and the  $\tilde{\chi}^0-N$  axial-vector interaction is induced by the  $t$ -channel exchange of one  $Z$  boson, respectively. The interactions of the neutralino LSP with the  $Z$  and Higgs bosons are suppressed by the mixing of gauginos and Higgsino. The squark exchange also generates the twist-2 interaction. The twist-2 coupling is not suppressed even in a pure gaugino limit, however, the term is proportional to  $m_{\tilde{q}}^{-4}$  in the amplitude. Overall, if the other SUSY particles are much heavier than the LSP and the mixing of gaugino and Higgsino is small, the elastic scattering is suppressed, as mentioned in Introduction. The one-loop contributions to the  $\tilde{\chi}^0-N$  scattering will be discussed in the next section.

#### A. Spin-independent interaction

The spin-independent interaction for the  $\tilde{\chi}^0-N$  scattering arises from one Higgs boson or squark exchange at tree-level. Since the squark contribution is typically subdominant, we concentrate on the Higgs boson contribution.

The neutralino coupling with a quark  $q$ ,  $f_q$ , from the Higgs boson exchange is given as

$$f_q[H] = \frac{g_2^2}{4m_W} \left( \frac{c_{h\tilde{\chi}\tilde{\chi}} c_{hqq}}{m_{h^0}^2} + \frac{c_{H\tilde{\chi}\tilde{\chi}} c_{Hqq}}{m_{H^0}^2} \right), \quad (19)$$

where

$$c_{hdd} = -\frac{\sin\alpha}{\cos\beta}, \quad c_{Hdd} = \frac{\cos\alpha}{\cos\beta} \quad (20)$$

for down-type quarks and

$$c_{huu} = \frac{\cos\alpha}{\sin\beta}, \quad c_{Huu} = \frac{\sin\alpha}{\sin\beta} \quad (21)$$

for up-type quarks. Here,  $\alpha$  and  $\beta$  are the mixing angle of the neutral Higgs bosons and the vacuum mixing angle, and  $m_{h^0}$  and  $m_{H^0}$  are the light and heavy CP-even Higgs boson masses, respectively. The tree-level coupling constants of the neutralino LSP with the Higgs bosons,  $c_{h\tilde{\chi}\tilde{\chi}}$  and  $c_{H\tilde{\chi}\tilde{\chi}}$ , are

$$c_{h\tilde{\chi}\tilde{\chi}} = [(O_N)_{12}^* - (O_N)_{11}^* t_W] \times [-\sin\alpha(O_N)_{13}^* - \cos\alpha(O_N)_{14}^*], \quad (22)$$

$$c_{H\tilde{\chi}\tilde{\chi}} = [(O_N)_{12}^* - (O_N)_{11}^* t_W][\cos\alpha(O_N)_{13}^* - \sin\alpha(O_N)_{14}^*],$$

where  $(O_N)$  is the neutralino mixing matrix. Here  $t_W = \tan\theta_W$ ,  $c_W = \cos\theta_W$  and  $s_W = \sin\theta_W$  with  $\theta_W$  the Weinberg angle.

For the winolike LSP, the couplings  $c_{h\tilde{\chi}\tilde{\chi}}$  and  $c_{H\tilde{\chi}\tilde{\chi}}$  are given as

$$c_{h\tilde{\chi}\tilde{\chi}} \approx \frac{m_W}{M_2^2 - \mu^2} (M_2 + \mu \sin 2\beta),$$

$$c_{H\tilde{\chi}\tilde{\chi}} \approx -\frac{m_W}{M_2^2 - \mu^2} \mu \cos 2\beta. \quad (23)$$

Here, we assume  $|\mu| - M_2 \gg m_Z$  and  $\cos\alpha \sim \sin\beta$  and  $\sin\alpha \sim -\cos\beta$ . The latter corresponds to a limit of the heavy pseudoscalar Higgs boson mass ( $m_A$ ). The coupling constants are suppressed when  $|\mu| \gg M_Z$ , as expected. When  $\tan\beta$  is large, the heavy Higgs boson contribution may be dominant. While the LSP coupling with the light Higgs boson is suppressed by  $\sim m_{\tilde{\chi}^0} m_W / \mu^2$ , the coupling of strange quark to the heavy Higgs boson is enhanced proportional to  $\tan\beta$ .

For the Higgsino-like LSP, we get

$$c_{h\tilde{\chi}\tilde{\chi}} \approx \mp t_W^2 \frac{1}{2} \frac{m_W}{M_1 - |\mu|} (1 \pm \sin 2\beta)$$

$$\mp \frac{1}{2} \frac{m_W}{M_2 - |\mu|} (1 \pm \sin 2\beta), \quad (24)$$

$$c_{H\tilde{\chi}\tilde{\chi}} \approx t_W^2 \frac{1}{2} \frac{m_W}{M_1 - |\mu|} \cos 2\beta + \frac{1}{2} \frac{m_W}{M_2 - |\mu|} \cos 2\beta,$$

$\mu > 0$  ( $\mu < 0$ ). In this paper we take the LSP mass positive by an axial rotation of the LSP field. While the couplings

are suppressed by the gaugino masses, the suppression is moderate compared with the winolike LSP.

Assuming that the dominant contribution comes from the light Higgs boson exchange, the cross section for the spin-independent  $\tilde{\chi}^0$ - $p$  scattering is approximately given as follows,

$$\sigma_{\text{SI}} \sim 3 \times 10^{-43} \text{ cm}^2$$

$$\times \left( \frac{m_{h^0}}{115 \text{ GeV}} \right)^{-4} \left( \frac{\mu^2}{100 \text{ GeV} \times M_2} \right)^{-2}$$

$$\times \left( 1 + \frac{\mu}{M_2} \sin 2\beta \right)^2 \quad (25)$$

for the winolike LSP, and

$$\sigma_{\text{SI}} \sim 1 \times 10^{-43} \text{ cm}^2 \times \left( \frac{m_{h^0}}{115 \text{ GeV}} \right)^{-4} \left( \frac{M_2}{100 \text{ GeV}} \right)^{-2} \quad (26)$$

for the Higgsino-like LSP. Here we assume  $M_2 \ll \mu$  ( $\mu \ll M_2 = M_1$ ) for the winolike (Higgsino-like) neutralino LSP. Note that the cross section for the winolike LSP is suppressed by  $\mu^{-4}$ . For  $\mu = 1$  TeV and  $M_2 = 100$  GeV and  $\tan\beta = 10$ , the cross section reduces down to  $10^{-46}$  to  $10^{-47}$   $\text{cm}^2$ . It is found in Ref. [20] that the wino- and Higgsino-like LSP masses are about (1–2) TeV and the spin-independent cross sections are  $10^{-(44-48)}$   $\text{cm}^2$ , when imposing the thermal relic density constraint.

## B. Spin-dependent interaction

The spin-dependent interaction of the LSP arises from one  $Z$  boson or one squark exchange. We ignore the squark contribution here, again. The tree-level contribution to  $d_q$  in Eq. (2) from the  $Z$  boson exchange is represented as

$$d_q = \frac{g_2^2}{8m_W^2} T_3^q c_{Z\tilde{\chi}\tilde{\chi}}, \quad (27)$$

where  $T_3^q$  is for the isospin of a quark  $q$ . The tree-level LSP coupling to  $Z$  boson,  $c_{Z\tilde{\chi}\tilde{\chi}}$ , is

$$c_{Z\tilde{\chi}\tilde{\chi}} = |(O_N)_{13}|^2 - |(O_N)_{14}|^2. \quad (28)$$

For the winolike LSP,  $c_{Z\tilde{\chi}\tilde{\chi}}$  becomes

$$c_{Z\tilde{\chi}\tilde{\chi}} \approx \frac{m_W^2}{M_2^2 - \mu^2} \cos 2\beta. \quad (29)$$

On the other hand, in a limit of the Higgsino-like LSP,

$$c_{Z\tilde{\chi}\tilde{\chi}} \approx \mp \frac{1}{2} \left( t_W^2 \frac{m_W^2}{M_1 \mu} + \frac{m_W^2}{M_2 \mu} \right) \cos 2\beta + O\left(\frac{\mu}{M_1}, \frac{\mu}{M_2}\right), \quad (30)$$

for  $\mu > 0$  ( $\mu < 0$ ). Thus, again, the coupling for the winolike LSP is more suppressed than that for the Higgsino-like one.

Using Eqs. (29) and (30), the spin-dependent cross section with proton is approximately given as

$$\sigma_{\text{SD}} \sim 2 \times 10^{-38} \text{ cm}^2 \times \left( \frac{\mu}{100 \text{ GeV}} \right)^{-4} \cos^2 2\beta \quad (31)$$

for the winolike LSP, and

$$\sigma_{\text{SD}} \sim 8 \times 10^{-39} \text{ cm}^2 \times \left( \frac{M_2}{100 \text{ GeV}} \right)^{-2} \left( \frac{\mu}{100 \text{ GeV}} \right)^{-2} \cos^2 2\beta \quad (32)$$

for the Higgsino-like LSP. Here, we take  $M_1 = M_2$ . The analysis in Ref. [20] shows that the spin-dependent cross sections are  $10^{-(41-45)} \text{ cm}^2$  for the wino- and Higgsino-like LSPs when imposing the thermal relic density constraint.

#### IV. ELASTIC SCATTERING INDUCED BY ONE-LOOP EFFECTIVE ACTION

In the previous section, we discuss that the interactions responsible for the  $\tilde{\chi}^0-N$  scattering are suppressed by the gaugino-Higgsino mixing or squark masses at tree-level. However, this is not true for the radiative corrections to the effective interactions if the dark matter is wino- or Higgsino-like, because of the mass degeneracy between the LSP and its  $SU(2)$  partner. In this section, we derive radiative corrections to the effective interactions in Eqs. (2), (3), and (11), and it is found that some of them are only suppressed by the weak gauge boson mass at most.

We first discuss the anomalous Higgs boson vertices of the neutralino and the box diagram contributions involving the  $W$  bosons to the effective interactions for the case of the winolike LSP. The numerical result will be shown in the next section. For the case of the Higgsino-like LSP, we present the explicit formula for the radiative corrections in the Appendix.

The gauge interactions of the winolike neutralino and chargino are

$$\begin{aligned} \mathcal{L}_{\text{int}} = & -\frac{e}{s_W} (\bar{\tilde{\chi}}^0 \gamma^\mu \tilde{\chi}^- W_\mu^\dagger + h.c.) + e \frac{c_W}{s_W} \bar{\tilde{\chi}}^- \gamma^\mu \tilde{\chi}^- Z_\mu \\ & + e \bar{\tilde{\chi}}^- \gamma^\mu \tilde{\chi}^- A_\mu. \end{aligned} \quad (33)$$

Here we ignore the mixings of the neutralinos and charginos for simplicity. These interactions induce the anomalous Higgs boson vertices and the box diagrams which contribute to the  $\tilde{\chi}^0-N$  scattering. The Feynman diagrams are given in Fig. 1. The radiative correction to the  $Z$  boson vertex is not induced at the one-loop level due to the Majorana nature of the LSP.

We start from the radiative correction to the the Higgs boson vertices. The tree-level contribution is given in Eq. (22). The radiative corrections to  $c_{h\tilde{\chi}\tilde{\chi}}$  and  $c_{H\tilde{\chi}\tilde{\chi}}$  at one-loop level are expressed as

$$\delta c_{h\tilde{\chi}\tilde{\chi}} = \frac{\alpha_2}{4\pi} \sin(\alpha - \beta) [F_H^{(0)}(x_W) + \delta_C F_H^{(1)}(x_W)], \quad (34)$$

$$\delta c_{H\tilde{\chi}\tilde{\chi}} = -\frac{\alpha_2}{4\pi} \cos(\alpha - \beta) [F_H^{(0)}(x_W) + \delta_C F_H^{(1)}(x_W)], \quad (35)$$

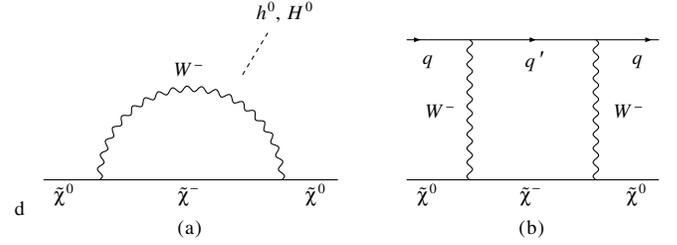


FIG. 1. (a) Diagram contributing to the anomalous Higgs boson vertices of the neutralino and (b) box diagram contributing to the  $\tilde{\chi}^0-N$  scattering in the case of the winolike LSP.

respectively. The correction to  $f_q[H](= f_q^{\text{tree}}[H] + \delta f_q[H])$  is therefore

$$\delta f_q[H] = \frac{g^2}{4m_W} \left( \frac{\delta c_{h\tilde{\chi}\tilde{\chi}} c_{hq} q}{m_{h^0}^2} + \frac{\delta c_{H\tilde{\chi}\tilde{\chi}} c_{Hq} q}{m_{H^0}^2} \right). \quad (36)$$

Here,  $x_W = m_W^2/m_{\tilde{\chi}^0}^2$  and  $\bar{b}_W = \sqrt{1 - x_W/4}$ . We expand the radiative corrections by  $\delta_C \equiv (m_{\tilde{\chi}^-} - m_{\tilde{\chi}^0})/m_{\tilde{\chi}^0}$ . When the gaugino and Higgsino masses are comparable to the weak gauge boson masses, this expansion of  $\delta_C$  is invalid, however, the one-loop contributions to the cross section are negligible. We are interested in a case where the gaugino or Higgsino mass is much larger than the weak gauge boson masses and the LSP is close to the Higgsino or wino weak eigenstate. This expansion is justified in the case.

The mass functions are

$$\begin{aligned} F_H^{(0)}(x) = & \frac{2}{\bar{b}_W} [2 + x(2 - x)] \tan^{-1} \left( \frac{2\bar{b}_W}{\sqrt{x}} \right) \\ & - 2\sqrt{x} [2 - x \log(x)], \end{aligned} \quad (37)$$

$$F_H^{(1)}(x) = \frac{6}{\bar{b}_W^3} \tan^{-1} \left( \frac{2\bar{b}_W}{\sqrt{x}} \right) - \frac{2}{\bar{b}_W^2} \frac{1}{\sqrt{x}} (2 + x). \quad (38)$$

Since  $F_H^{(0)}(x)$  becomes  $2\pi$  in a limit of  $x \rightarrow 0$ , the correction does not vanish in the heavy neutralino limit. When the pseudoscalar Higgs boson mass is large,  $\delta c_{h\tilde{\chi}\tilde{\chi}}$  is maximum and  $\delta c_{H\tilde{\chi}\tilde{\chi}}$  is vanishing.

The effective interactions in Eqs. (2) and (11) receive the radiative correction from the the box diagrams [Fig. 1(b)] at the renormalization scale  $\bar{\mu} \simeq m_W$  of the form,

$$\begin{aligned} \delta \mathcal{L}_q^{\text{eff}}[\text{box}] = & \delta d_q[\text{box}] \bar{\tilde{\chi}}^0 \gamma^\mu \gamma_5 \tilde{\chi}^0 \bar{q} \gamma_\mu \gamma_5 q \\ & + \delta f_q[\text{box}] m_q \bar{\tilde{\chi}}^0 \tilde{\chi}^0 \bar{q} q \\ & + \delta f'_q[\text{box}] \bar{\tilde{\chi}}^0 \tilde{\chi}^0 \bar{q} i \not{q} \\ & + \frac{\delta g_q^{(1)}[\text{box}]}{m_{\tilde{\chi}^0}} \bar{\tilde{\chi}}^0 i \partial^\mu \gamma^\nu \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^q \\ & + \frac{\delta g_q^{(2)}[\text{box}]}{m_{\tilde{\chi}^0}^2} \bar{\tilde{\chi}}^0 (i \partial^\mu) (i \partial^\nu) \tilde{\chi}^0 \mathcal{O}_{\mu\nu}^q, \end{aligned} \quad (39)$$

where

$$\begin{aligned}
\delta d_q[\text{box}] &= \frac{\alpha_2^2}{m_W^2} F_{\text{AV}}(x_W), \\
\delta f_q[\text{box}] &= \frac{\alpha_2^2}{m_W^3} F_{\text{S1}}(x_W), \\
\delta f'_q[\text{box}] &= \frac{\alpha_2^2}{m_W^3} F_{\text{S2}}(x_W), \\
\delta g_q^{(1)}[\text{box}] &= \frac{\alpha_2^2}{m_W^3} F_{\text{T1}}(x_W), \\
\delta g_q^{(2)}[\text{box}] &= \frac{\alpha_2^2}{m_W^3} F_{\text{T2}}(x_W).
\end{aligned} \tag{40}$$

In our calculation of the radiative corrections, we ignore  $O(m_q^2)$ , and expand the loop integrals up to order of  $p$  and  $\delta_C$ , where  $p$  is the quark external momentum. In this approximation the loop integrals are expressed analytically by  $B$  functions and its derivatives [21]. This procedure is not justified for box diagrams with the external or internal top quarks since it is heavier than the weak gauge bosons, however, the radiative corrections should be suppressed by the top quark mass. Thus, they are subdominant compared with the other lighter quark ones.

The loop function can be expanded as  $F_I(x) = F_I^{(0)}(x) + \delta_C F_I^{(1)}(x)$  ( $I = \text{AV}, \text{S1}, \text{S2}, \text{T1},$  and  $\text{T2}$ ), and they are given as follows;

$$\begin{aligned}
F_{\text{AV}}^{(0)}(x) &= \frac{1}{24\bar{b}_W} \sqrt{x}(8-x-x^2) \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) - \frac{1}{24} x[2 - (3+x) \log(x)], \\
F_{\text{AV}}^{(1)}(x) &= \frac{1}{4\bar{b}_W^3} \sqrt{x} \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) - \frac{1}{2\bar{b}_W^2}, \\
F_{\text{S1}}^{(0)}(x) &= F_{\text{S1}}^{(1)}(x) = 0, \\
F_{\text{S2}}^{(0)}(x) &= -\frac{\bar{b}_W}{24} (2+x^2) \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) - \frac{1}{96} \sqrt{x}[1 - 2x - x(2-x) \log(x)], \\
F_{\text{S2}}^{(1)}(x) &= -\frac{1}{24\bar{b}_W} (1-x)^2 \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) + \frac{1}{24\sqrt{x}} [2(6-x) + 6 \log(2\delta_C) - (3-x^2) \log(x)], \\
F_{\text{T1}}^{(0)}(x) &= \frac{1}{6} \bar{b}_W (2+x^2) \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) + \frac{1}{24} \sqrt{x}[1 - 2x - x(2-x) \log(x)], \\
F_{\text{T1}}^{(1)}(x) &= \frac{1}{6\bar{b}_W} (1-2x+x^2) \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) + \frac{1}{6\sqrt{x}} [2x + 6 \log(2\delta_C) - (3+x^2) \log(x)], \\
F_{\text{T2}}^{(0)}(x) &= \frac{1}{4\bar{b}_W} x(2-4x+x^2) \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) - \frac{1}{4} \sqrt{x}[1 - 2x - x(2-x) \log(x)], \\
F_{\text{T2}}^{(1)}(x) &= \frac{1}{24\bar{b}_W^3} (16 + 30x - 30x^2 + 5x^3) \tan^{-1}\left(\frac{2\bar{b}_W}{\sqrt{x}}\right) - \frac{1}{24\bar{b}_W^2} \sqrt{x}[28 - 10x - 5x(4-x) \log(x)].
\end{aligned} \tag{42}$$

Note that  $F_{\text{S1}}(x)$  is zero due to the chiral nature of  $W\bar{q}q$  vertex. For the Higgsino-like LSP, it does not vanish because the  $Z$  boson couples with both  $q_L$  and  $q_R$ . See the Appendix.

The functions  $F_{\text{S2}}^{(0)}(x)$  and  $F_{\text{T1}}^{(0)}(x)$  are nonvanishing even if  $x$  approaches to 0 (or  $m_{\tilde{\chi}^0}$  is increased), and they become  $-\pi/24$  and  $\pi/6$ , respectively. Thus, in addition to the correction to the Higgs boson vertices in Eq. (43), those to the scalar and twist two operators induced by the  $W$  boson loops do not vanish in a heavy LSP limit, and they contribute to the spin-independent cross section. On the other hand, the spin-dependent cross section depends on  $F_{\text{AV}}(x)$ , which is suppressed in the heavy LSP limit as  $\propto m_W/m_{\tilde{\chi}^0}$ .

The functions,  $F_{\text{H}}^{(1)}(x)$ ,  $F_{\text{S2}}^{(1)}(x)$  and  $F_{\text{T1}}^{(1)}(x)$ , are proportional to  $1/\sqrt{x}$  for small  $x$ . This does not cause a problem.

The mass difference  $\delta_C$  is proportional to  $x_W^{3/2}$  at tree-level, and it becomes proportional to  $\alpha_2 x_W^{1/2}$  due to the radiative correction when the LSP mass is much larger than the  $W$  boson mass. Thus,  $\delta_C/\sqrt{x_W} \sim x_W$  or  $\sim \alpha_2$ . Therefore the perturbation by  $\delta_C$  is not broken.

## V. NUMERICAL RESULTS

In the previous sections, we derived the formulas for the one-loop corrections to the effective interactions. We now calculate the cross section following the procedure in Sec. II.

First, we discuss the cross section for the spin-independent  $\tilde{\chi}^0$ - $p$  scattering induced by the gauge-loop diagrams presented in the previous section, assuming the

tree-level contribution to the cross section is negligible. The total cross section involving both the tree and one-loop level contributions is more sensitive to the MSSM model parameters, and it is discussed later.

For the winolike LSP, the spin-independent interaction induced by the  $W$  boson at one-loop level is approximately represented by

$$\begin{aligned} \delta f_N = & -\frac{\pi\alpha_2^2 m_N}{m_W m_{h^0}^2} \left( \sum_{q=u,d,s} \frac{f_{Tq}}{2} + \frac{f_{TG}}{9} \right) \\ & -\frac{\pi\alpha_2^2 m_N}{m_W^3} \left( \sum_{q=u,d,s} \frac{f_{Tq}}{24} + a_{\text{eff}} \frac{f_{TG}}{324} \right) \\ & + \sum_{q=u,d,s,c} \frac{\pi\alpha_2^2 m_N}{8m_W^3} [q(2, m_W^2) + \bar{q}(2, m_W^2)]. \quad (43) \end{aligned}$$

Here, we take the pseudoscalar Higgs boson and the LSP to be much heavier than the  $W$  boson mass. The first term comes from the correction to the light Higgs boson vertex, proportional to  $\delta c_{h\tilde{\chi}\tilde{\chi}}$ . The second one is from those to the scalar operators and it is given by  $\delta f_q$  and  $\delta f'_q$  in Eq. (39). The third (last) one represents those to the twist-2 operators, which are proportional to  $\delta g_q^{(1)}$  and  $\delta g_q^{(2)}$  in Eq. (39). Here, we include the top and bottom quark contributions to the Higgs vertex correction, while those to the scalar and twist-2 operators, which are suppressed by the top quark mass, are ignored. As discussed before,  $a_{\text{eff}}$  is a phenomenological parameter for  $\langle N | \bar{Q} i \not{D} Q | N \rangle$ . Equation (43) is numerically given as (in the order of the terms in Eq. (43))

$$\begin{aligned} \delta f_N = & -5.8 \times 10^{-10} - (5.3 + 1.6a_{\text{eff}}) \times 10^{-11} \\ & + 3.9 \times 10^{-10} (\text{GeV}^{-2}) \\ = & -2.6 \times 10^{-10} (\text{GeV}^{-2}), (a_{\text{eff}} = 1). \quad (44) \end{aligned}$$

Here,  $m_{h^0} = 115$  GeV and the parameters in Table I for the hadronic matrix elements are used. We assume  $a_{\text{eff}} = 1$  in the second line. We found the contributions from the terms with different spin structure cancel each other. This value for  $\delta f_N$  corresponds to  $\sigma_{\text{SI}} \approx 3.0 \times 10^{-47}$  cm<sup>2</sup>, if the tree-level contribution is negligible. The correction is still large enough to alter the total cross section close to the proposed sensitivities for the future experiments ( $10^{-(45-46)}$  cm<sup>2</sup>). Here, we present the spin-independent interaction of the neutralino with proton, however, the value for that with neutron is almost the same as it since the numerical difference comes only from the Higgs exchange contributions.

In Fig. 3, we show the one-loop induced cross section for the spin-independent  $\tilde{\chi}^0$ - $p$  scattering as a function of  $f_{Ts}$  for  $a_{\text{eff}} = 1.0$  and 2.0. The tree-level contribution is assumed to be negligible. The parameters  $f_{Ts}$  and  $a_{\text{eff}}$  have large theoretical uncertainties. The precise determination of  $a_{\text{eff}}$  requires calculation of two-loop diagrams, such as in Fig. 2. Here, we take  $m_{\tilde{\chi}^0} = 1600$  GeV and the param-

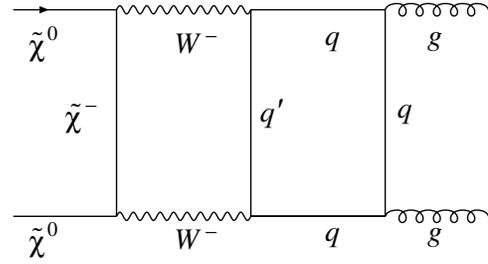


FIG. 2. A two-loop diagram contributing to  $GG\bar{\chi}^0\chi^0$ .

eters in Table I except for  $f_{Ts}$ . Also, we use FeynHiggs [22] in order to calculate the Higgs boson masses and mixing angle. In this figure, we use  $\tan\beta = 10$ ,  $m_A = 1000$  GeV,  $m_{\text{stop}} = 2000$  GeV, and  $A_{\text{top}} = 0$ , which lead to  $m_{h^0} = 116$  GeV. The cross section is very sensitive to  $f_{Ts}$  since the correction to the Higgs boson vertex is one of the dominant corrections. When  $f_{Ts}$  is smaller than 0.1, the cross section is significantly suppressed. Since the box diagram corrections to the scalar operators are numerically small, the  $a_{\text{eff}}$  dependence is relatively small.

In Fig. 4 the one-loop induced spin-independent cross section is presented as a function of  $\tan\beta$  and  $m_A$ . Here, we take  $m_{\tilde{\chi}_1^0} = 1600$  GeV,  $m_{\text{stop}} = 2000$  GeV,  $A_{\text{top}} = 0$ ,  $a_{\text{eff}} = 1$ , and the parameter in Table I. The  $\tan\beta$  dependence is moderate. The one-loop induced cross section rises as large as  $10^{-46}$  cm<sup>2</sup> when the pseudoscalar Higgs boson is light.

In the case of the Higgsino-like LSP, the one-loop induced cross section for the spin-independent scattering

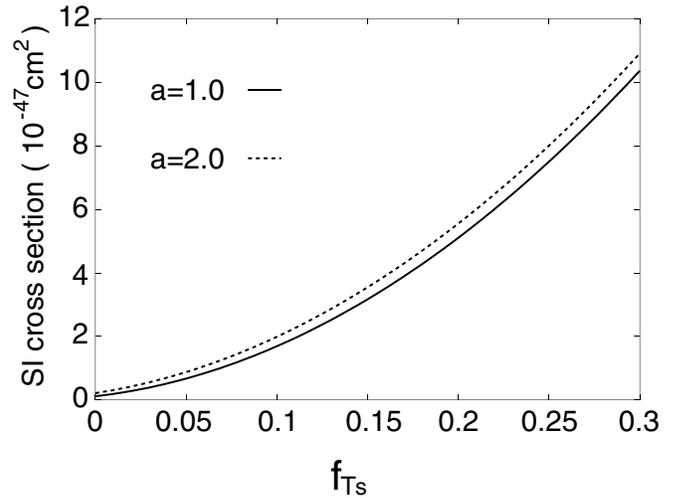


FIG. 3. One-loop induced cross section for spin-independent  $\tilde{\chi}^0$ - $p$  scattering as a function of  $f_{Ts}$  for  $a_{\text{eff}} = 1.0$  and 2.0 in the case of the winolike LSP. We assume that the tree-level contribution is negligible and the process is induced by the one-loop diagrams. Here, we take  $m_{\tilde{\chi}^0} = 1600$  GeV and the parameter in Table I except for  $f_{Ts}$ . Also, we use  $\tan\beta = 10$ ,  $m_A = 1000$  GeV,  $m_{\text{stop}} = 2000$  GeV, and  $A_{\text{top}} = 0$ , which determine the light Higgs mass and the mixing angle.

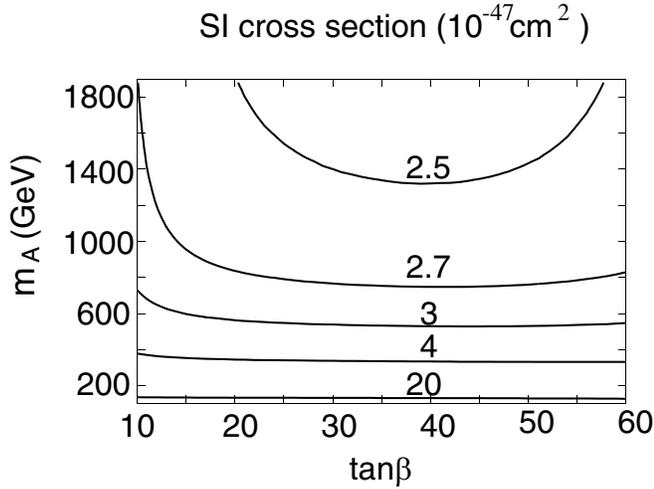


FIG. 4. One-loop induced cross section for spin-independent  $\tilde{\chi}^0-p$  scattering is presented as a function of  $\tan\beta$  and  $m_A$ . The winolike LSP is assumed. Here, we take  $m_{\tilde{\chi}^0} = 1600$  GeV,  $m_{\text{stop}} = 2000$  GeV,  $A_{\text{top}}=0$ ,  $a_{\text{eff}} = 1$ , and the parameter in Table. I.

is smaller than that in the winolike case since the  $SU(2)$  charge is smaller. When the pseudoscalar Higgs boson and the LSP are much heavier than the  $W$  boson mass,  $\delta f_N$  is given as

$$\begin{aligned} \delta f_N &= 4.3 \times 10^{-11} - (5.1 + 0.1 a_{\text{eff}}) \times 10^{-11} \\ &\quad + 6.0 \times 10^{-11} \\ &= 5.1 \times 10^{-11}, (a_{\text{eff}} = 1). \end{aligned} \quad (45)$$

Here, we take  $m_{h^0} = 115$  GeV, again. The order of the terms in the first line is the same as Eq. (44). This value for  $\delta f_N$  corresponds to  $\sigma_{\text{SI}} \simeq 1.1 \times 10^{-48}$  cm<sup>2</sup> when the tree-level contribution is negligible.

Now we discuss the gauge-loop correction in the total cross section in the general MSSM parameter space. In above paragraphs we showed the cross section induced by the one-loop diagrams alone, assuming the tree-level cross section is negligible. When the tree-level amplitude dominates, the one-loop correction to the cross section is approximately expressed as  $\sim 2\sqrt{\sigma_{1\text{-loop}}/\sigma_{\text{tree}}}$ , where  $\sigma_{\text{tree}}$

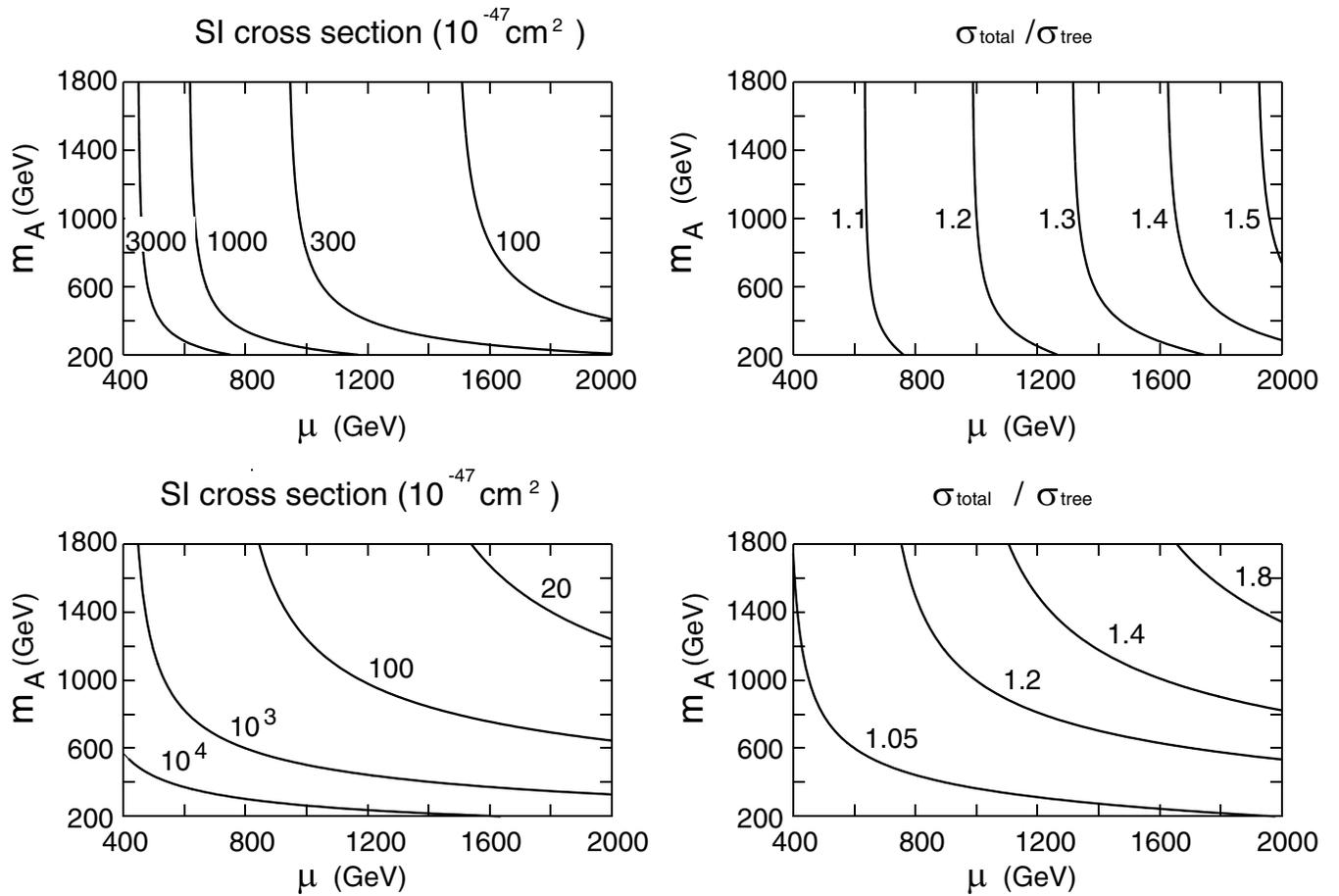


FIG. 5. Total cross section for spin-independent  $\tilde{\chi}^0-p$  scattering,  $\sigma_{\text{total}}$ , (left) for the winolike LSP in the  $(\mu, m_A)$  plane. The cross section is depicted in the unit of  $10^{-47}$  cm<sup>2</sup>. Here,  $M_2 = 200$  GeV,  $M_1/M_2 = 11t_W^2$ ,  $m_{\tilde{t}} = 2$  TeV and  $\tan\beta = 4(40)$  for upper (lower) plot.  $\sigma_{\text{total}}$  includes the gauge-loop contributions while the squark exchange contribution at tree-level is ignored. In the right, we give the ratio  $\sigma_{\text{total}}/\sigma_{\text{tree}}$  for the same parameters.  $\sigma_{\text{tree}}$  is the cross section at tree-level.

and  $\sigma_{1\text{-loop}}$  are the tree-level and one-loop induced cross sections, respectively. When  $\sigma_{1\text{-loop}}/\sigma_{\text{tree}} \sim 1/100$ , the correction to the total cross section is about 20%.

In Fig. 5 we show the cross section involving the tree-level and one-loop contributions,  $\sigma_{\text{total}}$ , and  $\sigma_{\text{total}}/\sigma_{\text{tree}}$  for the spin-independent  $\tilde{\chi}^0-p$  scattering as functions of  $\mu$  and  $M_A$ . Here, we assume the winolike LSP and take  $m_{\tilde{\chi}^0} \sim 200$  GeV. The upper two figures are for  $\tan\beta = 4$  and the lower ones are for  $\tan\beta = 40$ . See the caption for the other input parameters. It is found that the radiative correction is about 50% in the plots, where  $\sigma_{\text{total}}$  is above  $10^{-45}$  cm<sup>2</sup>. The radiative correction is relatively significant in larger  $\tan\beta$ , since the coupling of the LSP with the light Higgs boson at tree-level is more suppressed in the case as discussed in Section 3.1. In Fig. 6, we show the case where  $m_{\tilde{\chi}^0} \sim 2$  TeV. It is found that  $\sigma_{\text{total}}$  is less than  $10^{-45}$  cm<sup>2</sup>, however it is dominated by the gauge-loop correction in the plots, since the LSP is very close to the pure wino state.

In Figs. 7 and 8, we show the effect of the radiative correction for the Higgsino-like LSP. The radiative correction is less than  $\sim 10\%$  for  $m_{\tilde{\chi}^0} \sim 200$  GeV even when  $M_A$  and  $M_1$  are as heavy as 2 TeV. Compared with the winolike case, the suppression of the tree-level Higgs boson vertex by the gaugino-Higgsino mixing is more moderate while

the radiative correction by the gauge-loop diagrams is smaller. The correction is negative (positive) for  $\mu > 0$  ( $< 0$ ), relatively to the tree-level contribution. For  $m_{\tilde{\chi}^0} \sim 2$  TeV, the one-loop contribution cancels the tree-level contribution to reduce the total cross section less than  $10^{-50}$  cm<sup>2</sup> in the figure.

Finally, we discuss the spin-dependent cross section. As discussed in the previous section, the one-loop contribution to the process is suppressed by  $m_W/m_{\tilde{\chi}^0}$  in the amplitude, contrary to the spin-independent cross section. In Fig. 9 the one-loop induced cross section for the spin-dependent  $\tilde{\chi}^0-p$  scattering is given as a function of the LSP mass under the assumption of the winolike LSP. In this figure we assume again that the tree-level contribution is negligible. The asymptotic behavior for the cross section is

$$\begin{aligned} \sigma_{\text{SD}} &\simeq \frac{\pi\alpha_2^4 m_N^2}{3m_{\tilde{\chi}^0}^2 m_W^2} \left( \sum_{q=u,d,s} d_q \right) \\ &\simeq 1.2 \times 10^{-43} \text{cm}^2 \times \left( \frac{m_{\tilde{\chi}^0}}{100 \text{ GeV}} \right)^{-2}. \end{aligned} \quad (46)$$

Thus, from Eq. (31), it is found that the one-loop correction becomes significant for  $m_{\tilde{\chi}^0} \simeq 100$  GeV when  $\mu \gtrsim 1$  TeV. For  $|\mu| \gg M_2$ , the one-loop contribution is constructive to

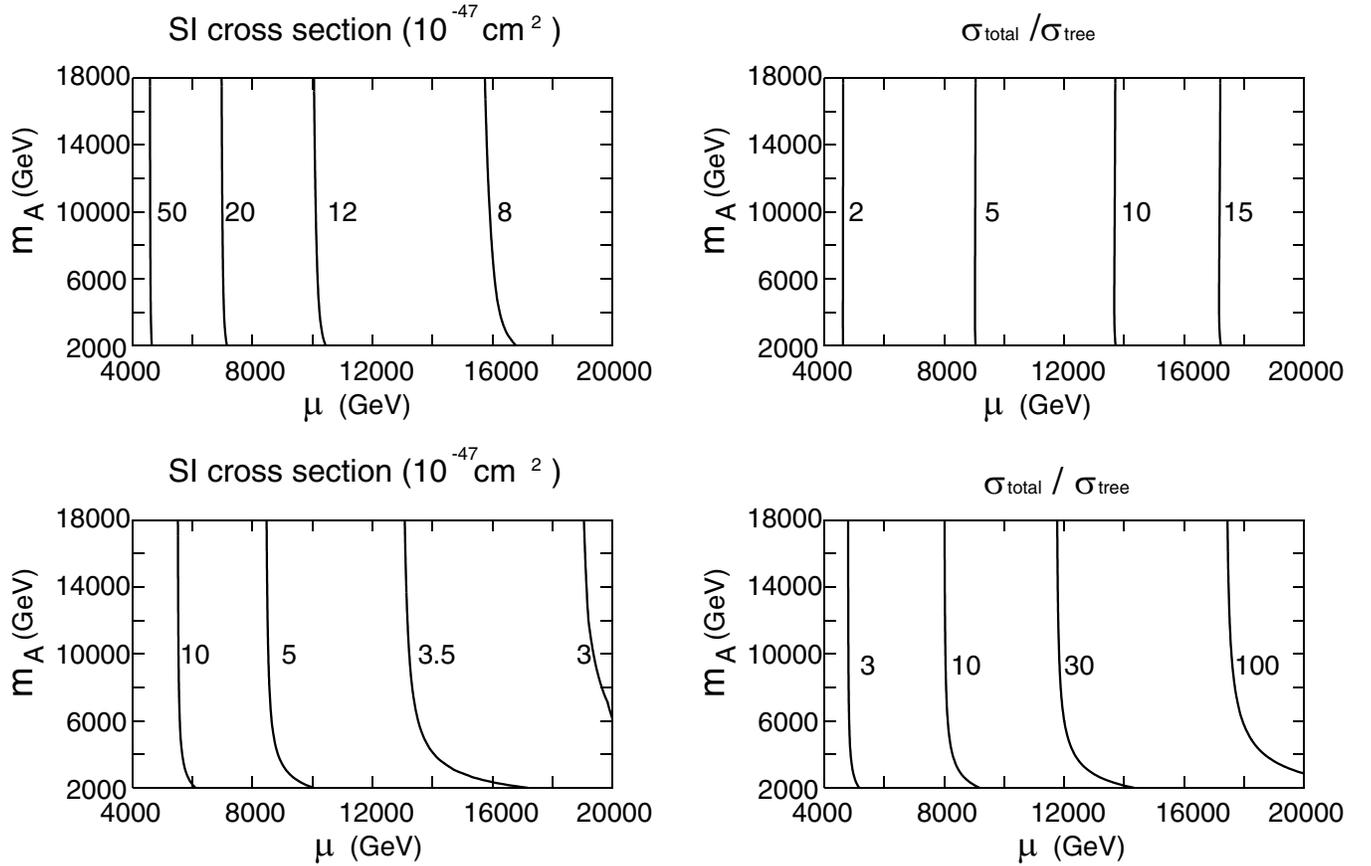


FIG. 6. Same as Fig. 5, but  $M_2 = 2$  TeV,  $m_i = 20$  TeV.

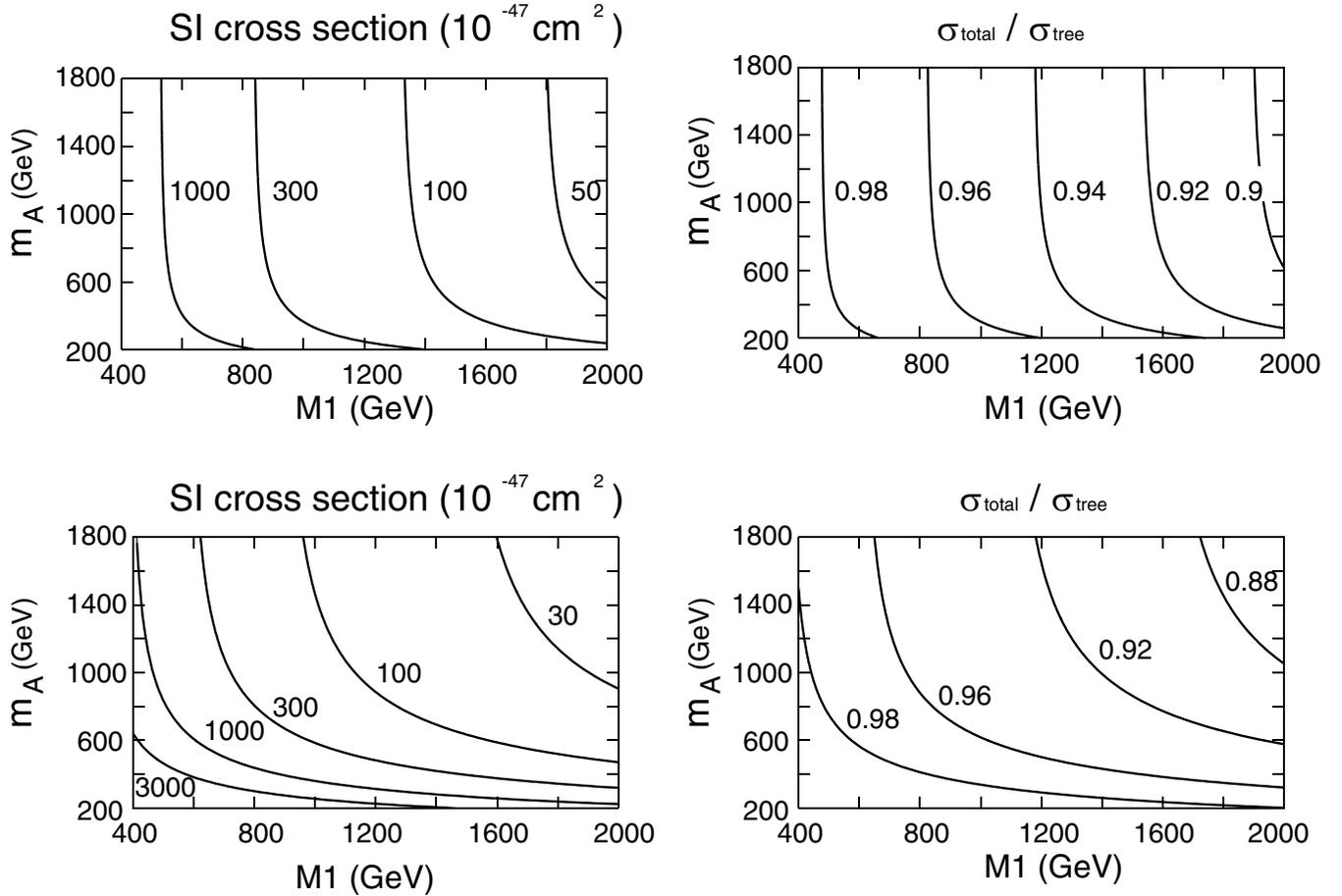


FIG. 7. Total cross section for spin-independent  $\tilde{\chi}^0-p$  scattering,  $\sigma_{\text{total}}$  (left) for the Higgsino-like LSP in the  $(\mu, m_A)$  plane. The cross section is depicted in the unit of  $10^{-47} \text{ cm}^2$ . Here,  $\mu = 200 \text{ GeV}$ .  $M_1/M_2 = 5t_W^2/3$ ,  $m_{\tilde{t}} = 2 \text{ TeV}$  and  $\tan\beta = 4(40)$  for upper (lower) figure, respectively. In the right, we give the ratio  $\sigma_{\text{total}}/\sigma_{\text{tree}}$ .

the tree-level one for the spin-dependent  $\tilde{\chi}^0-p$  scattering while it is deconstructive for the  $\tilde{\chi}^0-n$  scattering.

For the Higgsino-like LSP, the one-loop induced spin-dependent cross section is 1 order of magnitude smaller than that for the winolike LSP. It behaves as

$$\sigma_{\text{SD}} \approx 1.5 \times 10^{-44} \text{ cm}^2 \times \left( \frac{m_{\tilde{\chi}^0}}{100 \text{ GeV}} \right)^{-2} \quad (47)$$

in a large  $m_{\tilde{\chi}^0}$  limit. The tree-level contribution does not suffer from significant suppression by the Higgsino-gaugino mixing as in Eq. (32), and the one-loop correction is negligible as far as the gaugino masses are smaller than 10 TeV. For  $M_1, M_2 \gg |\mu|$ , the one-loop contribution is deconstructive (constructive) to the tree-level one for the spin-dependent  $\tilde{\chi}^0-p$  ( $\tilde{\chi}^0-n$ ) scattering.

## VI. CONCLUSION AND DISCUSSION

In this paper, we studied one-loop correction to the neutralino-nucleon scattering processes for the wino- and Higgsino-like LSPs. The  $\tilde{\chi}^0-N$  scattering is relevant to the detection rate of the neutralino dark matter. The scattering occurs dominantly through the exchange of the Higgs or Z

boson at tree-level, which is suppressed by the gaugino-Higgsino mixing. This is especially the case when the gauginos or Higgsino mass is large.

On the other hand, the scattering cross section receives the one-loop contribution from the processes involving the weak boson exchange. These processes are not suppressed neither by the small mixing angle nor by heavy SUSY mass scale, therefore it could be significant part of the one-loop corrected cross section in the limits of the wino- and Higgsino-like LSPs. The spin-independent cross section for the winolike (Higgsino-like) LSP receives the sizable one-loop correction, when the Higgsino (gaugino) mass is heavier than about 1 TeV and the spin-independent cross section is smaller than about  $10^{-45} \text{ cm}^2$  ( $10^{-46} \text{ cm}^2$ ).

There has been impressive progress on experimental techniques for the dark matter searches. The on-going experiments are now sensitive enough to exclude some of the MSSM parameters which may not be accessible otherwise. Many proposals aim to push the sensitivities further. The better determination on the  $\tilde{\chi}^0-N$  cross section is also needed to understand the local dark matter density and its velocity distribution. Although we calculated only a part of

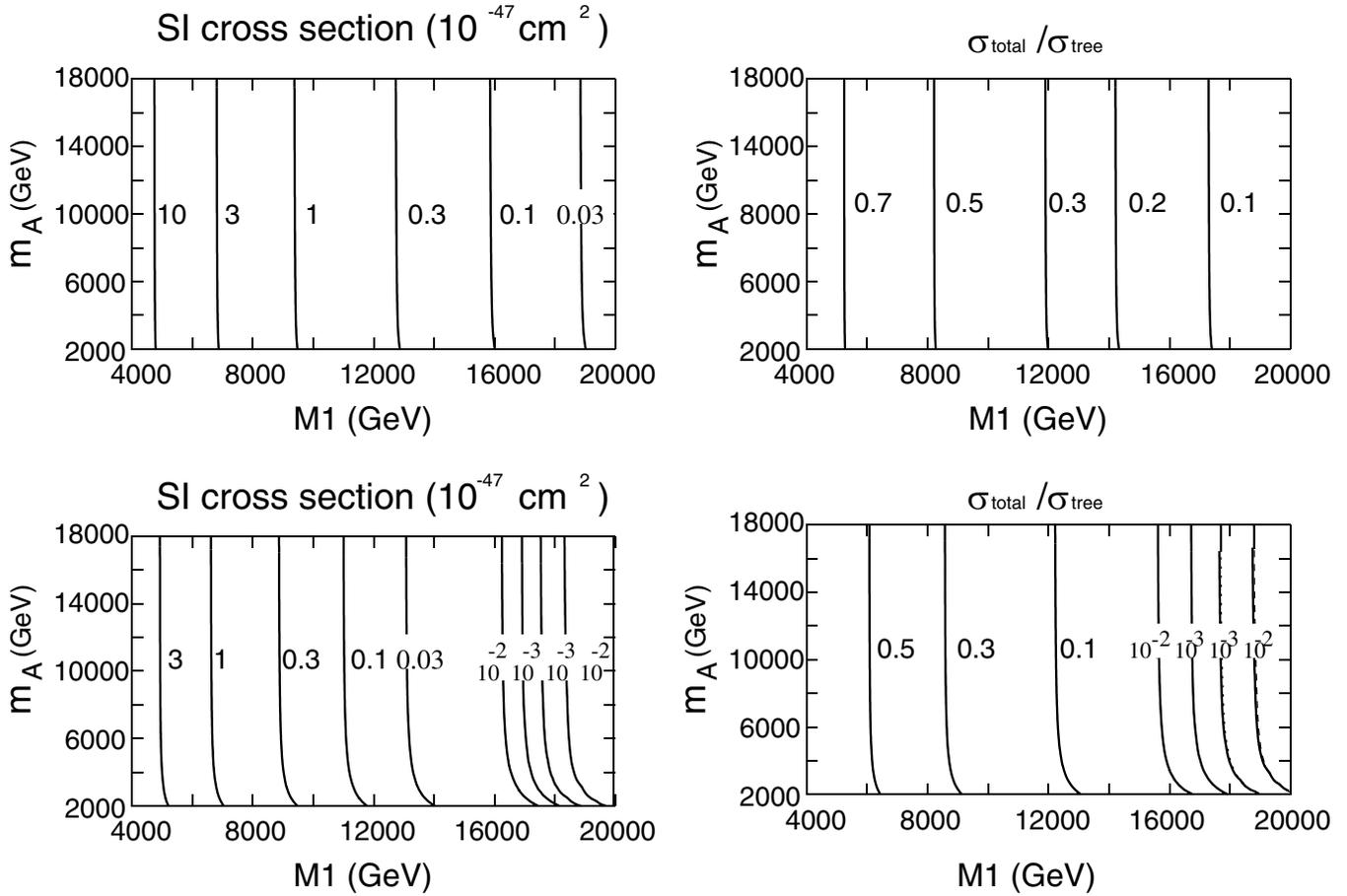


FIG. 8. Same as Fig. 7, but  $\mu = 2$  TeV and  $m_{\tilde{t}} = 20$  TeV.

the correction of the scattering cross section, further investigation on the loop effect might be needed in future.

**ACKNOWLEDGMENTS**

This work is supported in part by the Grant-in-Aid for Science Research, Ministry of Education, Science and Culture, Japan (No. 15540255, No. 13135207 and No. 14046225 for J.H. and No. 14540260 and No. 14046210 for M. M. N.).

**APPENDIX: HIGGSINO-LIKE LSP**

In this Appendix, we present the radiative corrections to the effective interactions for the LSP scattering with nucleon when the LSP is Higgsino-like. The Higgsino-like LSP accompanies the chargino and the second-lightest neutralino, whose masses are degenerate with that of the LSP. Furthermore, the Higgsino-like LSP has an interaction with the Z boson.

The one-loop corrections to the LSP couplings with the Higgs bosons are given as

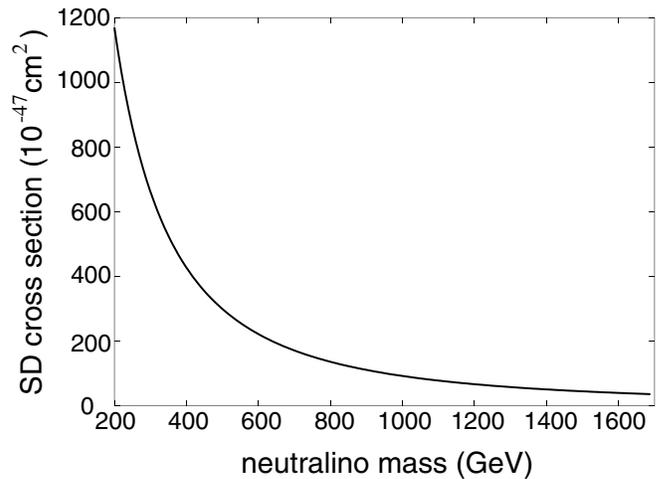


FIG. 9. One-loop induced cross section for the spin-dependent  $\tilde{\chi}^0-p$  scattering in the case of the winolike neutralino DM. Here, we assume that the tree-level contribution is negligible.

$$\begin{aligned} \delta c_{h\tilde{\chi}\tilde{\chi}}[\tilde{\chi}^-] &= \frac{1}{4} \frac{\alpha_2}{4\pi} \sin(\alpha - \beta) [F_H^{(0)}(x_W) + \delta_C F_H^{(1)}(x_W)] \\ &\quad - \frac{1}{4} \frac{\alpha_2}{4\pi c_W^2} \sin(\alpha - \beta) [F_H^{(0)}(x_Z) \\ &\quad + \delta_N F_H^{(1)}(x_Z)], \end{aligned} \quad (\text{A1})$$

$$\begin{aligned} \delta c_{H\tilde{\chi}\tilde{\chi}}[\tilde{\chi}^-] &= -\frac{1}{4} \frac{\alpha_2}{4\pi} \cos(\alpha - \beta) [F_H^{(0)}(x_W) \\ &\quad + \delta_C F_H^{(1)}(x_W)] + \frac{1}{4} \frac{\alpha_2}{4\pi c_W^2} \cos(\alpha - \beta) \\ &\quad \times [F_H^{(0)}(x_Z) + \delta_N F_H^{(1)}(x_Z)]. \end{aligned} \quad (\text{A2})$$

The contributions from the box diagrams to the coefficients in the effective Lagrangian (39) are approximated as

$$\begin{aligned} \delta d_q[\text{box}] &= \frac{1}{4} \frac{\alpha_2^2}{m_W^2} [F_{AV}^{(0)}(x_W) + \delta_C F_{AV}^{(1)}(x_W)] \\ &\quad - \frac{1}{2} (L_q^2 + R_q^2) \frac{\alpha_2^2}{c_W^2 m_W^2} \\ &\quad \times [F_{AV}^{(0)}(x_Z) + \delta_N F_{AV}^{(1)}(x_Z)], \end{aligned} \quad (\text{A3})$$

and

$$\delta f_q[\text{box}] = (L_q R_q) \frac{\alpha_2^2}{c_W m_W^3} [F_{S1}^{(0)}(x_Z) + \delta_N F_{S1}^{(1)}(x_Z)], \quad (\text{A4})$$

$$\begin{aligned} \delta f'_q[\text{box}] &= \frac{1}{4} \frac{\alpha_2^2}{m_W^3} [F_{S2}^{(0)}(x_W) + \delta_C F_{S2}^{(1)}(x_W)] \\ &\quad - \frac{1}{2} (L_q^2 + R_q^2) \frac{\alpha_2^2}{c_W m_W^3} \\ &\quad \times [F_{S2}^{(0)}(x_Z) + \delta_N F_{S2}^{(1)}(x_Z)], \end{aligned} \quad (\text{A5})$$

$$\begin{aligned} \delta g_q^{(I)} &= \frac{1}{4} \frac{\alpha_2^2}{m_W^3} [F_{TI}^{(0)}(x_W) + \delta_C F_{TI}^{(1)}(x_W)] \\ &\quad - \frac{1}{2} (L_q^2 + R_q^2) \frac{\alpha_2^2}{c_W m_W^3} [F_{TI}^{(0)}(x_Z) + \delta_N F_{TI}^{(1)}(x_Z)], \end{aligned} \quad (\text{I} = 1, 2). \quad (\text{A6})$$

Here,  $L_q = T_3^q - Q_q s_W^2$  and  $R_q = -Q_q s_W^2$  for a quark  $q$ . The  $Z$  boson contributions appear in the above equations. In the equations,  $x_Z = m_Z^2/m_{\tilde{\chi}^0}^2$ ,  $\bar{b}_Z = \sqrt{1 - x_Z/4}$ , and  $\delta_N \equiv (m_{\tilde{\chi}_2^0} - m_{\tilde{\chi}^0})/m_{\tilde{\chi}^0}$  with  $m_{\tilde{\chi}_2^0}$  the heavier neutralino mass. All mass functions are the same as in Eqs. (42) except for  $F_{S1}^{(0)}(x)$  and  $F_{S1}^{(1)}(x)$ ,

$$\begin{aligned} F_{S1}^{(0)}(x) &= \frac{1}{4\bar{b}_Z} [4 - x(2 - x)] \tan^{-1} \left( \frac{2\bar{b}_Z}{\sqrt{x}} \right) \\ &\quad + \frac{1}{4} \sqrt{x} [2 - x \log(x)], \end{aligned} \quad (\text{A7})$$

$$\begin{aligned} F_{S1}^{(1)}(x) &= \frac{3}{4\bar{b}_Z^2} (2 - x) \tan^{-1} \left( \frac{2\bar{b}_Z}{\sqrt{x}} \right) - \frac{1}{4\bar{b}_Z^2} \\ &\quad \times \frac{1}{\sqrt{x}} [4 - 4x - 2(4 - x) \log(2\delta_N) \\ &\quad + (4 - x) \log(x)]. \end{aligned} \quad (\text{A8})$$

- 
- [1] D.N. Spergel *et al.*, *Astrophys. J. Suppl. Ser.* **148**, 175 (2003).  
[2] C.L. Bennett *et al.*, *Astrophys. J. Suppl. Ser.* **148**, 1 (2003).  
[3] L. Randall and R. Sundrum, *Nucl. Phys. B* **557**, 79 (1999).  
[4] G. F. Giudice, M. A. Luty, H. Murayama, and R. Rattazzi, *J. High Energy Phys.* 12 (1998) 027.  
[5] A. Brignole, L. E. Ibanez, and C. Munoz, *Nucl. Phys. B* **422**, 125 (1994); **436**, 747(E) (1995).  
[6] T. Gherghetta, G. F. Giudice, and J. D. Wells, *Nucl. Phys. B* **559**, 27 (1999).  
[7] T. Moroi and L. Randall, *Nucl. Phys. B* **570**, 455 (2000).  
[8] N. Arkani-Hamed and S. Dimopoulos, hep-th/0405159; N. Arkani-Hamed, S. Dimopoulos, G.F. Giudice, and A. Romanino, hep-ph/0409232.  
[9] R. Bernabei *et al.*, *Riv. Nuovo Cimento* **26N1**, 1 (2003).  
[10] CDMS Collaboration, D. S. Akerib *et al.*, *Phys. Rev. Lett.*, **93** 211301 (2004).  
[11] L. Bergstrom and P. Ullio, *Nucl. Phys. B* **504**, 27 (1997); Z. Bern, P. Gondolo, and M. Perelstein, *Phys. Lett. B* **411**, 86 (1997).  
[12] J. Hisano, S. Matsumoto, and M. M. Nojiri, *Phys. Rev. D* **67**, 075014 (2003); *Phys. Rev. Lett.* **92**, 031303 (2004).  
[13] M. Drees and M. Nojiri, *Phys. Rev. D* **48**, 3483 (1993).  
[14] G. Jungman, M. Kamionkowski, and K. Griest, *Phys. Rep.* **267**, 195 (1996).  
[15] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. Nadolsky, and W. K. Tung, *J. High Energy Phys.* 07 (2002) 012.  
[16] European Muon Collaboration, J. Ashman *et al.*, *Nucl. Phys. B* **328**, 1 (1989).  
[17] H. Y. Cheng, *Phys. Lett. B* **219**, 347 (1989).  
[18] R. L. Jaffe and A. Manohar, *Nucl. Phys. B* **337**, 509

- (1990).
- [19] J. R. Ellis and R. A. Flores, Phys. Lett. B **263**, 259 (1991).
- [20] S. Profumo and C. E. Yaguna, Phys. Rev. D **70**, 095004 (2004).
- [21] G. 't Hooft and M. J. G. Veltman, Nucl. Phys. B **153**, 365 (1979).
- [22] S. Heinemeyer, W. Hollik, and G. Weiglein, Comput. Phys. Commun. **124**, 76 (2000).